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Ultravisible warped model from flavor triviality and improved naturalness

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A warped extra-dimensional model, where the standard model Yukawa hierarchy is set by UV physics, is shown to have a sweet spot of parameters with improved experimental visibility and possibly naturalness. Upon marginalizing over all the model parameters, a Kaluza-Klein scale of 2.1 TeV can be obtained at 2σ (95.4% C.L.) without conflicting with electroweak precision measurements. Fitting all relevant parameters simultaneously can relax this bound to 1.7 TeV. In this bulk version of the Rattazzi-Zaffaroni shining model, flavor violation is also highly suppressed, yielding a bound of 2.4 TeV. Nontrivial flavor physics at the LHC in the form of flavor gauge bosons is predicted. The model is also characterized by a depletion of the third-generation couplings—as predicted by the general minimal flavor violation framework—which can be tested via flavor precision measurements. In particular, sizable CP violation in $\Delta B = 2$ transitions can be obtained, and there is a natural region where B_s mixing is predicted to be larger than B_d mixing, as favored by recent Tevatron data. Unlike other proposals, the new contributions are not linked to Higgs or any scalar exchange processes.

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I. INTRODUCTION

Plunging the standard model (SM) in a warped extra-dimension provides new perspectives on understanding electroweak symmetry breaking (EWSB), offering a new way to solve the gauge hierarchy problem [1]. The Randall-Sundrum (RS) class of models also offers a simple way to address the SM flavor puzzle by localizing the SM fermions away from the Higgs vacuum expectation value (VEV) with $\mathcal{O}(1)$ parameters [2,3], which is referred to as the anarchic approach. In addition, the anarchic setup protects against large flavor and CP violation via the so-called RS-GIM mechanism [4–6]. Yet, a residual little CP problem, in the form of too large contributions to ϵ_K [7–11] and electric dipole moments [4,5,12], remains. (Some more RS flavor issues can be found in e.g. [13–22].) Furthermore, this framework calls for improvement on naturalness since a fine-tuning of worse than $\mathcal{O}(10\%)$ [9,23–25] of the electroweak (EW) scale is required to comply with EW precision tests (EWPTs) [26,27]. In the best of all known RS models, including a custodial symmetry for $Z \to b\bar{b}$, the lore is that this pushes the Kaluza-Klein (KK) scale above 3 TeV (below we argue that these numbers may be too optimistic).

It has been known for some time that changing the position of the light fermions, thus giving up on the virtues of the anarchic approach, may result in a better EW fit. In particular, if the profile of all the light fermions is close to being flat, a suppression of the Peskin-Takeuchi S parameter is obtained [26,28–33]. This would allow one to lower the KK-scale and possibly improve the naturalness of the model. It is interesting that such a fermion setup is consistent with imposing in the bulk the approximate SM flavor symmetries: $U(2)_Q \times U(2)_U \times U(2)_D \times U(3)_L \times U(3)_E$, where Q, U, D (L, E) correspond to the SM quark (lepton) doublet, up- and down-type quark (charged lepton) singlets, respectively.

In the following, we propose to give up on the warped extra-dimensional built-in mechanism for solving the flavor puzzle and the RS-GIM protection; after all, no experimental evidence implies that the flavor hierarchies arise from TeV scale physics, while, on the other hand, the hierarchy problem does inevitably point to it. We assume that the Yukawa hierarchy is set by some unknown physics on the UV brane, while both the bulk and the IR brane are invariant under the (now gauged) SM flavor symmetries.[1] Then the hierarchical five-dimensional (5D) fundamental Yukawa couplings are shined through the bulk by scalar flavon fields, thus realizing the approximate SM flavor symmetry structure.

Such a setup was first proposed by Rattazzi and Zaffaroni (RZ) [36], where the SM fields were localized on the IR brane as in the original RS1 model [1]. In this case, higher-dimensional operators, which generically contribute to EWPTs and flavor-changing neutral currents (FCNCs), can be suppressed, but only at the expense of a severe little hierarchy problem.

[1]The lepton symmetry can be also be broken down to products of $U(2)$. However, for simplicity we do not consider this possibility nor do we focus on lepton flavor violation, which is suppressed in our framework, or neutrino masses. Both issues are discussed in [34,35].

We show below that a bulk version of the RZ model, with a bulk Higgs [37], leads to a very exciting class of models, where improved agreement with EWPTs is obtained. We perform a global fit to the EW precision observables, evaluating the contributions to S, T, and $Z \to b\bar{b}$ at one-loop order (which are calculable in this model). In order to compare our model to the celebrated anarchic case, we also repeat the fit for this case. (However, for simplicity, we only consider one possible custodial assignment.) As a result of our analysis, we find a sweet spot in the parameter space of the bulk RZ model, which allows for a significantly lower KK-scale, such that it would be much easier to observe (or exclude [38]) at the LHC. Furthermore, we show that the inclusion of the one-loop contributions to the EW observables raises the KK-scale of the anarchic case. In addition, the fine-tuning associated with our model is ameliorated relative to the anarchic case.

The above scenario offers also some interesting perspective on flavor physics. First of all, the ϵ_K RS problem is solved, so that the bound from flavor is considerably weakened. Second, the model is characterized by a depletion of the third-generation couplings, as predicted by the general minimal flavor violation framework [39]. The model also yields sizable CP violation (CPV) in $\Delta B = 2$ transitions with, in particular, the possibility to obtain CPV contributions in B_s mixing larger than in B_d, as seems favored by the Tevatron data at present [40]. This is achieved without invoking Higgs or other scalar exchange processes [41–43]. Finally, since the bulk flavor symmetry is gauged, such that large breaking effects from quantum gravity are avoided, flavor gauge bosons are awaited around the TeV scale. Such states may be discovered at the LHC [44].

In short, the main differences between our study and previous ones are:

(i) We give a rational and an explicit model (a bulk RZ setup with some rough speculations on a possible extension to grand unification), where the light fermion profiles are roughly flat.

(ii) We choose a custodial representation for the leptons, which turns out to significantly improve the result of the global EW fit.

(iii) We emphasize, by calculating explicitly (and via 5D power counting), that in the bulk Higgs case the S parameter is one-loop finite. Furthermore, our estimation of the UV sensitive contribution to S, based on naive dimensional analysis (NDA), shows that they are subdominant (for related discussions, see [27,33,45]). Thus, the resulting value of S is dominated by finite contributions, and is under control.

(iv) We use updated input parameters for our EW fit taken from [46]. Additionally, an appropriate top 5D Yukawa value, matched to the top mass at the relevant scale, is used, and our 5D gauge couplings are matched at one loop.

(v) Our statistical treatment consists of two different analyses. In the first one we report a bound on the KK-scale upon marginalizing over all the other model parameters, while in the second we produce a bound when all the relevant parameters are combined in a multidimensional fit.

(vi) Finally, even though for simplicity we have not considered the case where the Higgs is a pseudo-Goldstone boson (PGB), we provide a rough speculation on the fine-tuning of PGB extensions that include the above improvements, which can be compared to other genuine PGB studies [9,23,25].

The rest of the paper is organized as follows. In Sec. II we describe the warped 5D setup and define our notation. Then, in Sec. III, the constraints from EWPTs on this class of models are presented, while in Sec. IV we elaborate on their flavor phenomenology. Finally, Sec. V gathers our conclusions and discusses prospects at the LHC.

II. THE MODEL

We work in a slice of AdS_5 space-time. The metric is $ds^2 = (kz)^{-2}(\eta_{\mu\nu}dx^\mu dx^\nu - dz^2)$ with $\eta_{\mu\nu} = \mathrm{diag}(+---)$ and a curvature scale $k \simeq 2.4 \times 10^{18}$ GeV, hence solving the hierarchy problem all the way up to the Planck scale. The slice is bounded by two branes at $z = R \sim k^{-1}$ and $z = R' \sim \mathrm{TeV}^{-1}$ usually referred to as the UV and IR branes, respectively. We impose a $SU(2)_L \times SU(2)R_L \times U(1)_X$ gauge symmetry in the bulk. For simplicity, in this study we assume that the Higgs field—$H \sim (2,2)_0$ under the $(L,R)_X$ custodial gauge group—is a bulk field with VEV $\langle H\rangle = v_5(z,\beta)/\sqrt{2}$, where $v_5(z,\beta) \simeq vR'/R^{3/2}\sqrt{2(1+\beta)}(z/R')^{2+\beta}$ and $v \simeq 246$ GeV [47]. The β parameter sets the VEV localization in the bulk, with $\beta = 0$ corresponding to gauge-Higgs unified models [23,48]. Eventually, the model should be lifted to one where the Higgs is realized as a pseudo-Goldstone boson [23,49], so that the quadratically divergent corrections to the Higgs mass are cut at the KK-scale.[2] Therefore, we choose $\beta = 0$ in the following, so we expect our conclusions to approximately hold also in models where the Higgs is a pseudo-Goldstone boson. We also gauge in the bulk the non-Abelian part of the SM flavor symmetry $SU(3)_Q \times SU(3)_U \times SU(3)_D \times SU(3)_L \times SU(3)_E$, such that all flavor-changing effects are controlled by the SM Yukawas, thus realizing the minimal flavor violation (MFV) ansatz [36,50,51]. The fermions are embedded as $Q \sim (2,2)_{2/3}$, $U \sim (1,1)_{2/3}$, $D \sim (1,3)_{2/3} \oplus (3,1)_{2/3}$ and $L \sim (2,2)_{-1}$, $E \sim (1,3)_0 \oplus (3,1)_0$, so they transform covariantly under the custodial parity [48].

The bulk gauge symmetry breaks down to the SM gauge group on the UV brane and still preserves a custodial $SU(2)_{L+R}$ after EWSB, so the T parameter is protected

[2]We leave this specific analysis to future work.

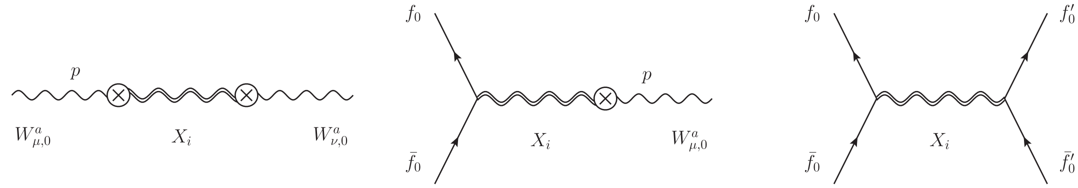

FIG. 1. Tree diagrams contributing at leading order to the EWPO. The double line denotes a sum over the various gauge KK-states, while the cross represents KK/zero-mode mixing from the Higgs VEV. W^a_0 are the SM zero modes with $a = 0, \ldots, 3$ and $W^0 \equiv B$ is the hypercharge gauge field.

from large bulk cutoff corrections. The breaking of the flavor group occurs only on the UV brane, and is shined toward the IR by some flavon scalar fields Φ, with VEV $\langle\Phi\rangle \propto Y_I$, where Y_I are the 5D Yukawa matrices ($I = U$, D, E). In contrast with most previous studies, we take the 5D Yukawas to display the hierarchy observed in 4D, which boils down to assuming that the latter are set by unspecified UV physics. The large top Yukawa implies a shift in the third-generation bulk masses, while the 5D bottom Yukawa is free to be taken either large or small. The latter can be regarded as the large or small $\tan\beta$ cases in two Higgs doublet models, such as supersymmetric theories. For simplicity we shall assume in the EW global fit that the 5D bottom Yukawa is small, and leave the implications of a large bottom Yukawa option to the flavor physics discussion in Sec. IV. This setup guarantees that at low energies the model belongs to the MFV framework [52–61], where harmless top Yukawa resummation is expected and may be observable in the future. Note that although taking a somewhat larger 5D bottom Yukawa (but still suppressed compared to the 5D top Yukawa) would not strongly affect the EWPTs, it would lead to a richer flavor phenomenology. In addition, flavor violation from the presence of flavor gauge bosons is also expected, but yet again, it is going to be subject to MFV protection [44]. In the following we discuss in more detail the EW and flavor sectors of our model and their phenomenological implications.

III. ELECTROWEAK PRECISION TESTS

Models of new physics for the EW scale are tightly constrained, at the *per mile* level, by the measurements at SLD and LEP, both at and above the Z pole [62,63], as well as by the Tevatron. In a large set of such models, the gauge sector observables, described by the so-called oblique parameters, capture most (if not all) of the constraints on new physics. Moreover, the large coupling of the top to the EWSB sector typically implies sizable nonoblique corrections for the third-generation quarks, notably to the $Z\bar{b}_L b_L$ coupling. The oblique parameters, along with the Z partial decay width into $b\bar{b}$, constitute a reduced set of EW precision observables (EWPOs) often sufficient to constrain RS models [24,26,27,64–67]. Indeed, whenever localized toward the UV brane, the (elementary) light fermions are barely sensitive to EWSB in the IR and induce negligible corrections to the EWPOs. In contrast, since the light fermions are more composite in our setup, additional nonoblique corrections are expected to be generated. This requires a more careful study of other observables, such as the hadronic Z decay width and observables sensitive to four-fermion operators in the lepton sector, like atomic parity violation (APV) in heavy nuclei. In such a highly nonuniversal new physics model, this implies that one must look at more than $\mathcal{O}(35)$ EWPOs in order to assess whether EWPTs are passed. In the following, we discuss in detail how such a fit is performed. Then, we report the resulting bounds on the KK-scale and estimate the fine-tuning in our model (as well as in the anarchic case) by computing the sensitivity of the new physics scale to the input parameters.

A. Global fit to electroweak observables

In order to properly include all possible correlations among the various observables, we perform a global fit to the EW precision data following the approach of [68,69]. To do so, we match the relevant dimension-six operators in the SM to our RS setup (see [65] for a review), including the most important, top (and eventually bottom) Yukawa enhanced, radiative corrections to the S and T parameters and the $Z\bar{b}_L b_L$ vertex. Radiative corrections to lighter fermion-gauge boson couplings and to four-fermion operators will be suppressed by smaller Yukawas. Note that when the bottom 5D Yukawa is $\mathcal{O}(1)$ or larger, additional loop contributions to the $Z\bar{b}_L b_L$ vertex involving neutral currents become important. We have not included such contributions and therefore will limit our analysis to a relatively small bottom Yukawa, where these radiative corrections are subdominant.[3]

We include the following RS contributions to the SM dimension-six operators. First of all, working at leading order in $(v/m_{\mathrm{KK}})^2 \ll 1$, the tree-level effects arise from exchange of KK-gauge bosons through the diagrams of Fig. 1. Additional tree-level contributions from the left-handed bottom sector (potentially, controlled by the top Yukawa coupling) are absent due to the custodial protection [48] (more generally, the left-handed down-type sector of the three generations enjoys a custodial protection). For

[3]In practice, this implies a hierarchy of ~4–5 between the 5D bottom Yukawa and the best fit value obtained for the top one.

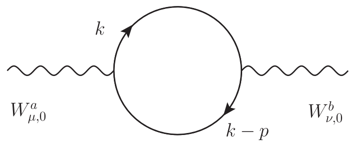

FIG. 2. Diagram contributing to the SM gauge boson propagators at one loop.

the up-type, as well as the right-handed down quark sectors, which are not protected by this symmetry, the effects are suppressed by the assumed hierarchical nature of the 5D Yukawa couplings (except for the top which is not, at present, experimentally constrained).

Furthermore, it is known that isospin breaking in the fermionic sector leads to sizable corrections to T at one loop [26]. This correction is often negative as a result of the choice of custodial representations, unless the singlet contribution dominates, in which case T can be positive at one loop [27,45]. On the other hand, the one-loop corrections to S tend to be positive and relatively small in RS for a reasonable range of parameters [27,45]. To prevent the appearance of a large S parameter at tree level and cancel the effect on the global fit of the small positive one-loop correction, we will focus on a region where the light fermions are almost flat [26]. Notice that since S is not protected by any symmetry, it could *a priori* be UV sensitive in 5D, whereas T is finite to all orders in perturbation thanks to the custodial symmetry. However, we show below that for a bulk Higgs the S parameter is one-loop finite. Thus the one-loop shifts in *both* S and T are calculable and dominated by the first KK-states. In practice, we include the first two KK-levels in the fit; higher KK-levels would yield at most a $(m^{(1)}_{\rm KK}/m^{(3)}_{\rm KK})^2 \sim 10\%$ correction, which we choose to neglect. Moreover, third-generation KK-quarks dominate the shift to the weak gauge boson two-point functions through the diagram of Fig. 2, while other contributions are suppressed by either gauge couplings or smaller 5D Yukawas.

We include the one-loop correction to the $Zb_L\bar{b}_L$ coupling as well. The dominant contribution is from KK-fermions through the diagrams of Fig. 3, involving the SM charged current. Finally, although such contributions are not present in the model under study, we report the impact on the fit of the corrections to the S and T parameters arising from SM loops with a pseudo-Goldstone Higgs [70,71]. We refer the reader to Appendix A for further details on both the tree and one-loop calculations.

B. UV Sensitivity of the S parameter

We start by deriving the 5D degree of divergence of various one-loop contributions to S using NDA. We match the various relevant diagrams onto the coefficient C_S of the 5D local operator, $B_{\mu\nu}W_L^{\mu\nu a}H^\dagger\sigma^a H$, that generates S in the 4D effective action via $S = 4\pi v^2 C_S/gg'$. Gauge and fermion contributions to this operator scale as $C_S^g \propto g_5^4$ and $C_S^Y \propto Y_t^2 g_5^2$, respectively. Recalling that the Yukawa coupling has the same mass dimension as the 5D gauge coupling for a bulk Higgs, $[Y_t] = [g_5] = -1/2$, power counting yields $C_S^{g,Y} \sim \Lambda_5^{-1}$, hence a finite contribution, where $\Lambda_5 \equiv N_{\rm KK} \times k$ is the 5D cutoff. Thus S is perfectly calculable at one loop, and is dominated by the KK-fermion contribution, provided $Y_t \gg g_5$.

The 5D top Yukawa grows fast in the UV and quickly becomes nonperturbative. A conservative approach usually requires $N_{\rm KK} \gtrsim 3$, so that the 5D construction makes sense as an effective theory; we choose $N_{\rm KK} = 3$ in the following. Assuming in addition that the KK-fermion coupling to a bulk Higgs is $\mathcal{O}(1)$ for $\beta = 0$, NDA yields a perturbativity upper bound on Y_t of

$$Y_t\sqrt{k} \le 4\pi/\sqrt{N_{\rm KK}} \simeq 7.3. \tag{1}$$

Higher loops, however, will still be divergent, as they involve more powers of Y_t and/or g_5. This introduces a UV cutoff sensitivity, even for a bulk Higgs, starting at the two-loop level. Nonetheless, we argue that the S parameter calculation is still under control. Indeed, as exemplified by the diagram shown in Fig. 4, the two-loop correction scales like Y_t^4 or $Y_t^2 g_5^2$, so its contribution to S diverges like $\log N_{\rm KK}$. In Fig. 4 we show only the Higgs as the internal line. As shown below, this is justified for the sweet spot parameters, for which $Y_t\sqrt{k} \sim 5$. Hence contributions from an exchange of KK-gauge bosons will be subdominant, since they are proportional to $g_5^2 k \sim 9$ (see Appendix A), leading to a $g_5^2/Y_t^2 \sim 36\%$ correction NDA then yields

$$S^{\rm NDA}_{\text{2-loop}} \simeq \frac{4\pi v^2}{m_{\rm KK}^2}\frac{N_c}{(16\pi^2)^2} Y_t^4 k^2 \log N_{\rm KK}, \tag{2}$$

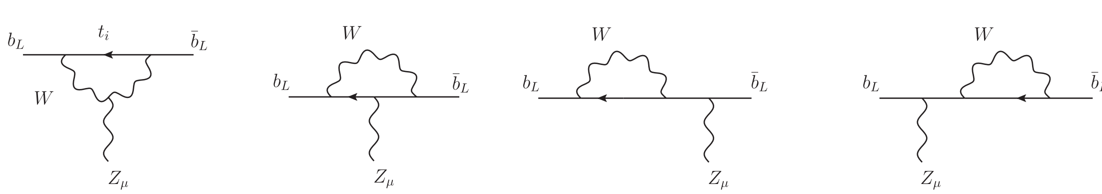

FIG. 3. One-loop diagrams contributing to $Zb_L\bar{b}_L$ in the unitary gauge. KK-modes of third-generation $Q = 2/3$ states and $W^\pm$ zero-mode are running in the loop.

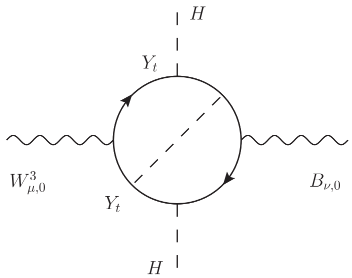

FIG. 4. Two-loop diagram relevant for matching onto the dimension-six operator generating the S parameter. A similar diagram with exchange of weak gauge boson is also present.

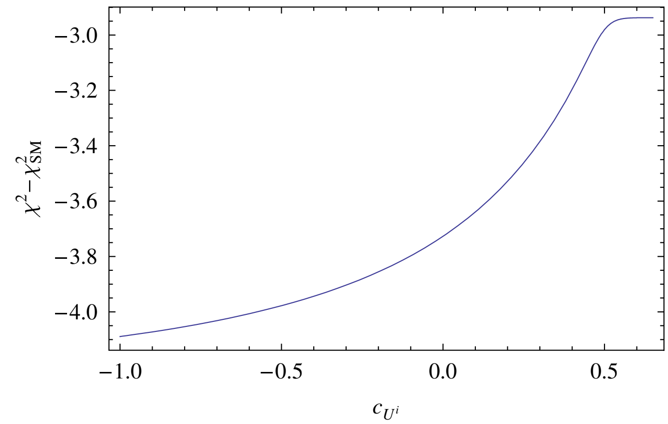

FIG. 5 (color online). $\chi^2 - \chi^2_{\rm SM}$ as a function of c_{U^i}, with the rest of the parameters fixed to the best fit values of Eq. (3). Besides being relatively insensitive to the localization of U^i, the χ^2 distribution flattens for $c_{U^i} < -0.5$. Here we take $m_H = 115$ GeV.

where we used the fact that KK-fermion coupling to a bulk Higgs is $\mathcal{O}(1)$ for $\beta = 0$ [18]. Thus, $S^{\rm NDA}_{2\text{-loop}}$ is suppressed by about $Y_t^2/16\pi^2 \log N_{\rm KK} \sim 20\%$ compared to the one-loop correction. Higher loops will be even more suppressed since, according to NDA, the expansion parameter is $Y_t^2\Lambda_5/16\pi^2$, which is smaller than 1 for a perturbative Yukawa. The S parameter is therefore under control in our setup.

C. Statistical analysis

We first count the parameters of interest in our model. Imposing the MFV ansatz, the bulk SM flavor symmetries receive large breaking only from the third-generation quark Yukawa couplings. Then, the $SU(3)3_{Q,U,D}$ flavor bulk symmetry is broken down to an approximate $SU(2)3_{Q,U,D}$ by the flavon VEVs, while the lepton flavor group is unbroken. Therefore, the whole set of effective operators in the SM is determined by 9 free parameters, which we choose to be the fermion bulk masses: c_{Q^3}, c_t, c_b, c_{Q^i}, c_{U^i}, and c_{D^i} ($i = 1, 2$, with universal first two generation masses) for the quark sector, c_L and c_E for the leptons (also taken to be family universal), and the KK-scale, $m_{\rm KK}$. The flavon VEVs are set by the SM fermion masses.

The global fit analysis proceeds as follows. First of all, a χ^2-distribution is constructed by comparing the experimental measurements to the theoretical predictions of the model; it is therefore a function of the new physics parameters: $\chi^2 = \chi^2(x)$, where, in our case, x collectively denotes the 9 parameters listed above. The most probable parameter values, $\bar{x}$, are then identified by minimizing the total χ^2 function with respect to the model parameters: $\chi^2(x = \bar{x}) \equiv \chi^2_{\rm min}$. Finally, we bound the parameters x to lie within confidence level regions around $\bar{x}$, whose size and shape are dictated by the χ^2 difference, $\Delta\chi^2(x) \equiv \chi^2(x) - \chi^2_{\rm min}$. The value of $\Delta\chi^2$ is fixed as a function of the chosen confidence level (C.L.) and the number of simultaneously constrained parameters.

For this analysis we assume a light Higgs and fix, for definiteness, its mass to $m_H = 115$ GeV. We find that the "best fit" parameters in the present scenario are

$$\begin{aligned} &m_{\rm KK} = 3.5\ \text{TeV}, \qquad c_t \simeq 0.47, \qquad c_b \simeq 0.6, \\ &c_{Q^i} \simeq 0.54, \qquad c_L \simeq 0.47, \qquad c_e \simeq 0.50, \end{aligned} \tag{3}$$

with a considerably lower sensitivity of the fit to c_{Q^3}, c_{U^i}, and c_{D^i} (at the minimum of the χ^2, we find $c_{Q^3} \simeq 0$ and $c_{D^i} \simeq 0.76$). As shown in Fig. 5, there is a preference for U^i to be composite, although the χ^2 does not depend strongly on c_{U^i} when U^i is sufficiently IR localized. The values given in Eq. (3) correspond to $c_{U^i} \simeq -0.5$, which we will use as a benchmark point. In addition, we impose the restriction $c_b \leq 0.6$ in order to ensure that the 5D bottom Yukawa coupling is sufficiently small, so that the $Q = -1/3$ states give a negligible contribution to $\delta g_{Z\bar{b}_L b_L}$ (we have not included such contributions; the required calculation can be extracted from [72]). Note, however, that the fit prefers a value at the allowed upper limit for c_b.

The goodness-of-fit of the above model is found to be $\chi^2_{\rm min}/\text{d.o.f.} = 217.3/223 \approx 0.97$. This can be compared to the goodness-of-fit of the SM with a Higgs mass $m_H = 90$ GeV (currently the best fit value): $\chi^2_{\rm SM}/\text{d.o.f.} = 219.9/232 \approx 0.95$. Thus, the agreement of this particular beyond the standard model scenario is quite comparable to the SM.[4] We note, however, that we did not fit the SM input parameters in Eq. (3), but rather fixed them to their best fit values in the absence of new physics [46]. We proceed now to set C.L. limits for models that deviate from Eq. (3).

[4]The net decrease in the total χ^2 with respect to the SM can be traced to $\sigma_{\rm had}$, R_e, R_μ, A_e, and to a number of LEP II cross sections. Conversely, we find a worse fit to the forward-backward asymmetry of the bottom, $A^{(0,b)}_{FB}$, and to a lesser extent to R_b. We also assume in our fit a Higgs mass $m_H = 115$ GeV, but this has a negligible impact on the total χ^2. For example, the SM with $m_H = 115$ GeV has $\chi^2 = 221.3$, which is only about 1.4 larger than the value given above.

In our scenario, $\Delta\chi^2$ is a function of 9 new input parameters, displaying a smooth decoupling limit, with approximately $\Delta\chi^2 \propto m_{\rm KK}^{-2}$. We choose to present bounds at the 95.4% (2σ) C.L. We describe two statistical treatments, of distinct physical relevance, to bound the KK-scale from EWPTs at a given confidence level. First of all, we derive a one d.o.f. bound on $m_{\rm KK}$ only, by marginalizing[5] over all the bulk masses and imposing $\Delta\chi^2 = 4.00$. In addition, we quote a bound on the KK-scale resulting from a simultaneous fit of the most relevant model parameters. We adopt the following simple criterion for assessing the relevance of a given input parameter: if $\partial \log\Delta\chi^2/\partial \log c_i > \mathcal{O}(1)$, then the parameter c_i is relevant in setting the C.L. We find that the logarithmic derivatives with respect to c_{Q^3}, c_{U^i} and c_{D^i} are much smaller than 1 (so we do not count them as d.o.f. for setting the C.L. limits), while the rest of the logarithmic derivatives are order 1 or larger. Therefore, the second bound on $m_{\rm KK}$ is obtained by assuming $9 - 3 = 6$ d.o.f., which translates into $\Delta\chi^2 = 12.8$ for 95.4% C.L.

We stress that these two bounds have a meaning of their own and contain complementary information. For the one d.o.f. analysis, we have that, statistically, 95.4% of the models show a KK-scale larger than the bound, *without any assumptions on all the other parameters*. Thus, we expect the one d.o.f. bound on the KK-scale to be rather conservative and of most relevance in terms of LHC discovery potential. On the other hand, the six d.o.f. analysis informs us on the possible correlations among the model parameters and, in particular, on the existence of less constrained directions in the parameter space. The presence of the latter could allow for a lower $m_{\rm KK}$, provided some other parameters deviate from their best fit values in a correlated way. As a result, however, we expect such a KK-scale to be statistically unlikely, although we have not tried to quantify this statement. Nevertheless, we think that the existence of such points in the parameter space are worth mentioning, for such correlation may be theoretically motivated and/or future experimental analyses may become sensitive to additional parameters, on top of the KK-scale.

D. EWPT global fit results

We report in this section the sweet spots found for the one and six d.o.f. statistical analyses defined above. These are done for both our flavor-triviality model, as well as for a (slight) variant of the conventional anarchic model.

[5]Assuming the parameters to be Gaussian distributed, marginalizing over the bulk masses boils down to setting them to the values that minimize the χ^2 as a function of the KK-scale: i.e., $c_i = c_i(m_{\rm KK})$, where the $c_i(m_{\rm KK})$'s satisfy a null gradient condition $\partial\chi^2/\partial c_i|_{m_{\rm KK}} = 0$. (See e.g. [73].)

1. Sweet spot in the flavor-triviality model

Assuming the hierarchical Yukawa ansatz, we find from the global fit a bound on the KK-scale of $m_{\rm KK} > 2.1$ TeV (95.4% C.L.) for the one d.o.f. analysis (i.e. $\chi^2 - \chi^2_{\rm min} = 4.00$). The corresponding sweet spot values for the bulk masses are

$$\begin{aligned} &c_{Q^3} \simeq 0.05, \quad c_t \simeq 0.47, \quad c_{Q^i} \simeq 0.51, \quad c_L \simeq 0.48, \\ &c_e \simeq 0.50, \quad c_b \simeq 0.6, \quad c_{U^i} \simeq -0.5, \quad c_{D^i} \simeq 0.77. \end{aligned} \tag{4}$$

If instead the SM is taken as the "best fit point" (i.e. $\chi^2 - \chi^2_{\rm SM} = 4.00$), the resulting bound is improved to 1.8 TeV (with some changes in the bulk masses). Performing a six d.o.f. analysis at 95.4% C.L. (i.e. $\chi^2 - \chi^2_{\rm min} = 12.8$) yields $m_{\rm KK} > 1.7$ TeV, with

$$\begin{aligned} &c_{Q^3} \simeq 0.02, \quad c_t \simeq 0.48, \quad c_{Q^i} \simeq 0.50, \quad c_L \simeq 0.48, \\ &c_e \simeq 0.50, \quad c_b \simeq 0.6, \quad c_{U^i} \simeq -0.18, \quad c_{D^i} \simeq 0.77. \end{aligned} \tag{5}$$

We illustrate in Fig. 6, as a function of the left-handed lepton localization parameter, c_L, the $\Delta\chi^2$ contributions which are most sensitive to this parameter; they are the Z pole observables (including b and c quark observables) and

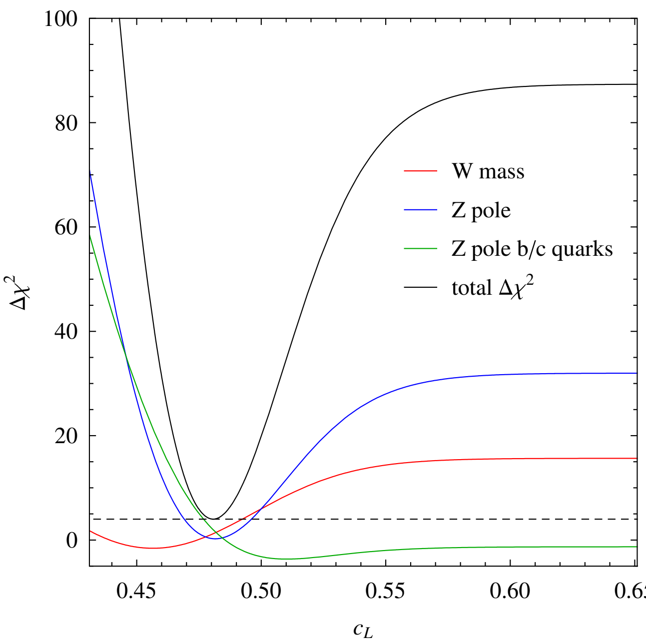

FIG. 6 (color online). Most important contributions to $\Delta\chi^2$ as a function of c_L, with $m_{\rm KK}$ and other bulk masses set to the sweet spot values of Eq. (4). The different curves, ordered according to their value on the right side of the plot from bottom to top, correspond to Z pole heavy quark observables including $R_{b,c}$, $A_{b,c}$, and $A^{FB}_{b,c}$ (green), W mass measurements (red), Z pole observables including total Z width, e^+e^- hadronic cross section and other leptonic observables (blue), and total $\Delta\chi^2$ (black), respectively.

TABLE I. Contributions to the χ^2 and $\Delta\chi^2$ for the sweet spot of Eq. (4); see [68].

	$\chi^2 - \chi^2_{\mathrm{SM}}$	$\chi^2_{\min} - \chi^2_{\mathrm{SM}}$	$\Delta\chi^2$
W mass	1.37	0.12	1.25
Z line shape and lepton A_{FB}	−2.47	−2.74	0.27
Z pole b and c quarks	5.30	3.44	1.86
s^2_W hadronic charge asymmetry	0.20	0.23	−0.03
Leptonic polarization asymmetries	−1.95	−2.18	0.23
Deep inelastic scattering	−0.13	−0.12	−0.01
Atomic parity violation	3.23	0.11	3.12
LEP2 hadronic cross section	−1.97	−0.97	−1.00
LEP2 muon pair	$<10^{-2}$	0.03	−0.03
LEP2 tau pair	−0.04	−0.03	−0.01
OPAL electron pair	−0.02	−0.02	$<10^{-2}$
L3 W pair	−0.17	−0.11	−0.06
Z pole s quark	0.07	0.09	−0.02
LEP2 $ee \to bb$	−3.22	−1.75	−1.47
LEP2 $ee \to cc$	−0.18	−0.08	−0.10
Total	0.02	−3.98	4

the W mass measurements. This shows that a low KK-scale is achieved for relatively flat light fermions, $c_L \simeq 0.48$, as expected from cancellation of an effective S parameter [see also Eq. (13) below]. We also report in Table I the contributions to the χ^2 and $\Delta\chi^2$ for the one d.o.f. sweet spot of Eq. (4).

2. Bounds on the semianarchic model

For the sake of comparing our setup to known anarchic models, and better assessing the benefits of the flavor-triviality scenario, we report the EW global results for a semianarchic model defined as follows. We set the first two quark generations and all the leptons to be elementary, $c_{Q^i} = c_{U^i} = c_{D^i} = c_L = c_E = 0.65$, thus allowing the corresponding 5D Yukawa couplings to be all of the same order (the fit is completely insensitive to the precise value of the c's, or to the fact that these are all the same, as long as they are UV localized). However, we require $c_b < 0.6$, so that the 5D bottom Yukawa coupling is suppressed. This restriction ensures that our loop contribution to $\delta g_{Z\bar{b}_L b_L}$ is reliable, as mentioned above. The results of relaxing this assumption, so that full anarchy can be achieved, will be presented elsewhere. The parameters determined in the fit are m_{KK}, c_{Q^3}, c_t, and c_b. Unlike what was found in the flavor-triviality model, under the anarchy assumption the minimum χ^2 is obtained when $m_{\mathrm{KK}} \to \infty$, hence it is the same as in the SM. The goodness-of-fit in this case is $\chi^2_{\mathrm{SM}}/\mathrm{d.o.f.} = 221.3/228 \approx 0.97$.

We then find a one d.o.f. bound ($\chi^2 - \chi^2_{\mathrm{SM}} = 4.00$) on the KK-scale of $m_{\mathrm{KK}} > 4.6$ TeV (95.4% C.L.) with the following sweet spot values:

$$c_{Q^3} \simeq 0.11, \qquad c_t \simeq 0.49, \qquad c_b \simeq 0.6. \tag{6}$$

For completeness, we also report that comparing the semianarchic scenario to the best fit point of Eq. (3), the bound is raised to $m_{\mathrm{KK}} \gtrsim 7$ TeV ($\chi^2 - \chi^2_{\min} = 4.00$).

In order to set a limit on m_{KK} by simultaneously fitting all the parameters, we note that m_{KK} and c_t are unequivocally relevant parameters (as defined in the previous subsection), while the logarithmic derivative of $\Delta\chi^2$ with respect to c_{Q^3} is much smaller than 1, and the one corresponding to c_b is of order 1. We therefore perform a $4 - 1 = 3$ d.o.f. analysis, corresponding to $\Delta\chi^2 = 8.02$ for 95.4% C.L. This yields $m_{\mathrm{KK}} > 3.9$ TeV with

$$c_t \simeq 0.49, \qquad c_b \simeq 0.6, \tag{7}$$

where we fixed $c_{Q^3} = 0.10$.

We end this subsection by emphasizing that in this work we explore the possibility that the fermions span $SU(2)_L \times SU(2)_R$ representations. However, the loop-level contributions to S, T, and $\delta g_{Z\bar{b}_L b_L}$ can be rather dependent on this assumption. For instance, when the third-generation fermions are assigned to $SO(5)$ representations, *à la* gauge-Higgs unification, one finds that the one-loop S parameter can be significantly smaller than for the $SU(2)_L \times SU(2)_R$ representations (this happens, e.g., in the scenario of Ref. [45]). This can affect the bounds for the anarchic scenario, which are controlled by the oblique parameters (plus $\delta g_{Z\bar{b}_L b_L}$). As an example, in the scenario discussed in [45], where the corresponding bound was found to be $m_{\mathrm{KK}} > 3.4$ TeV, an updated 3-parameter fit to the EW data leads to $m_{\mathrm{KK}} > 3$ TeV (both cases are compared to the SM as the best fit). This slight improvement is mainly due to the use of the most recent SM fit, that has moved in a favorable direction for these scenarios. Regarding the flavor-triviality case, we expect that the representation assignments for the third family are less crucial for the sweet spot, since the difference in the loop-level S parameter can be compensated to some extent by an effective tree-level contribution (with a slight readjustment of parameters). What is more important for the sweet spot is the custodial protection associated with the lepton representations, as emphasized in the introduction. With these caveats, we explore the degree of tuning involved, in the next subsection.

E. Fine-tuning estimates

The bulk RZ model displays a sweet spot of bulk masses where the KK-scale is significantly lower than for the most optimistic anarchic cases. Such a lower KK-scale would, in principle, reduce the fine-tuning associated with the Higgs mass in warped models. However, this result is a mere consequence of the approximate flatness of the SM light fermion wave functions. We therefore expect a large sensitivity of the KK-scale under corrections to the sweet spot values of the bulk masses, and a potentially larger new source of fine-tuning. We show that, even in the anarchic case, a sensitivity of this sort is actually also present; we shall estimate its size as well.

Because of the flavor symmetries, the only UV sensitive contributions are expected to be related to gauge interactions, which distinguish between different fermion representations. This would raise a legitimate question regarding the sweet spot: why are fields related to different SM representations located near each other, in particular, around $c_x \sim 0.5$? We do not have a sharp answer to this question. However, one could imagine embedding the above theory into some form of unification model (for an $SO(10)$ grand unified theory (GUT) see for example [74–80]), which would explain why the couplings are related to each other.[6] Finite radiative corrections to these quantities are proportional to the bulk masses themselves [81]. The radiative corrections to the bulk masses, which split the universal part of the fermions' wave functions, c_i, will be finite and suppressed by a loop factor of order $g_5^2 k/16\pi^2$ ($g_5^2 k \sim 9$). Therefore, a mass splitting of a few percent is expected. It is interesting that within the RS framework the radiative correction to the masses seems to vanish for flat fermions. However, we shall not pursue this possibility, as it goes beyond the scope of this project.

We now estimate the fine-tuning in the flavor triviality and semianarchic models. The fine-tuning is composed of two ingredients, namely, the sensitivity of the weak scale to the KK-scale, $FT_{m_{\rm KK}}$, and the sensitivity of the KK-scale itself to the bulk masses, FT_c, through the global EW fit. Strictly speaking, the former is under control only if the Higgs is a pseudo-Goldstone boson (i.e., if its mass is finite). Nonetheless, we believe a pragmatic approach[7] consists in estimating this fine-tuning as $FT_{m_{\rm KK}} \simeq (v/f_\pi)^2$, where $f_\pi^{-1} \approx R'$ (see e.g. [23]). The specific definition is not crucial (and different studies differ in their order 1 coefficients in any case [9,23,25]), but it does enable us to assess the difference in success between the anarchic and flavor triviality cases.

Regarding FT_c, one conventional procedure is to relate it to the logarithmic derivative at the sweet spot, $S_c \equiv \max_i|\partial \log m_{\rm KK}/\partial \log c_i|$ [82,83], and to interpret its inverse as a measure of the fine-tuning involved. However, since the sweet spot naturally resides in a local minimum of the parameter space, this derivative exactly vanishes. Instead, we find it convenient to use the (one-sided) finite difference analog, which gives a measure of the sensitivity of $m_{\rm KK}$ to c_i in a vicinity of the sweet spot. Since these one-sided finite derivatives can be different on both sides of the best fit point, especially for parameters which control the size of the top Yukawa coupling (e.g. c_t), we use an average of the two,

$$S_{c_i} \equiv \frac{1}{2}\left(\left| \frac{c_i}{\Delta c_i} \frac{\Delta m_{\rm KK}^+}{m_{\rm KK}} \right| + \left| \frac{c_i}{\Delta c_i} \frac{\Delta m_{\rm KK}^-}{m_{\rm KK}} \right| \right), \qquad (8)$$

where we choose $\Delta c_i = 0.03$, as motivated by the typical size of the radiative corrections. Here $\Delta m_{\rm KK}^{\pm} = m_{\rm KK}(c_i \pm \Delta c_i) - m_{\rm KK}(c_i)$ is the change of the KK-scale for a given Δc_i, with the other bulk masses fixed, that is necessary to keep $\Delta\chi^2$ fixed (so as to keep the success of the EW global fit at the same level). The final sensitivity should correspond to the largest value obtained by repeating the procedure for all the parameters in the model, $S_c = \max_i S_{c_i}$.

Using the definitions above, we find for the flavor-triviality (one, six) d.o.f. sweet spots the following measures of fine-tuning:

$$\begin{gathered} FT_{m_{\rm KK}} \simeq (8.4, 13)\%; \qquad S_{c_t}^{-1} \simeq (5.3, 10)\%; \\ S_{c_L}^{-1} \simeq (6.2, 8.9)\%; \qquad S_{c_E}^{-1} \simeq (8.1, 9.3)\%; \\ S_{c_{Q^i}}^{-1} \simeq (46, 42)\%, \end{gathered} \qquad (9)$$

[6]This would require various fermions to come from the same GUT multiplet, which would require a nonconventional approach to ensure proton longevity. Alternatively, one could impose a discrete symmetry that would correspond to invariance with respect to interchanging the different GUT multiplets.

[7]In the present model the Higgs mass parameter is quadratically sensitive to the cutoff scale, rather than to the KK-scale. Our intention here is to very roughly estimate the fine-tuning of the EW scale in PGB extensions that also incorporate the ingredients discussed in this paper. One should remember, however, that such extensions can contain additional correlations that may not allow a KK-scale as low as we have found above, or may contain additional sources of fine-tuning (see e.g. [25]). Nevertheless, we believe that the new ingredients highlighted here should help in relaxing the bound on $m_{\rm KK}$, and hence associated the fine-tuning in such models.

while there is essentially no sensitivity to the rest of the input parameters. If the SM is assumed to constitute the best fit, these numbers slightly improve,

$$\begin{aligned} FT_{m_{\rm KK}} &\simeq (11, 14)\%; \qquad S^{-1}_{c_t} \simeq (8.1, 10)\%; \\ S^{-1}_{c_L} &\simeq (8.2, 9.1)\%; \qquad S^{-1}_{c_E} \simeq (9.2, 9.4)\%; \\ & S^{-1}_{c_{Q^i}} \simeq (44, 42)\%. \end{aligned} \tag{10}$$

In contrast, for the semianarchic model with fermions in $SU(2)_L \times SU(2)_R$ representations, we find

$$\begin{aligned} \text{SM:}\ FT_{m_{\rm KK}} &\simeq (1.7, 2.3)\%; \quad S^{-1}_{c_t} \simeq (22, 20)\%, \\ \text{Best fit:}\ FT_{m_{\rm KK}} &\simeq (0.7, 1.7)\%; \quad S^{-1}_{c_t} \simeq (-, 22)\%, \end{aligned} \tag{11}$$

for (one, three) d.o.f., respectively.[8] We then see that indeed the fine-tuning of the weak scale is improved in our model. On the other hand, the sensitivity to the bulk masses is greater. We regard the sensitivity exhibited by S_c^{-1} as an indication of fine-tuning that, together with $FT_{m_{\rm KK}}$, determines the overall fine-tuning of the model.

Finally, it is interesting to analytically examine the sensitivity of the vertex corrections [see Eqs. (A4) and (A5)] to the localization of the leptons in our model, which is the main source for $S^{-1}_{c_{L,E}}$, as given above. The parametric dependence on $c_{L,E}$ and $m_{\rm KK}$ of the corresponding operators is

$$a^{t,s}_{hF} \propto \frac{(c_{L,E} - 1/2 + k)}{m^2_{\rm KK}}, \tag{12}$$

where $c_{L,E} - 1/2$ originated from the $\{++\}$ gauge KK-states [26] and $k \simeq 0.06$ effectively parametrizes the contribution of the $\{-+\}$ gauge KK-states. Using this expression, we can approximately evaluate $S^{-1}_{c_{L,E}}$ as

$$S^{-1}_{c_{L,E}} = \left(\frac{\partial \log m_{\rm KK}}{\partial \log c_{L,E}}\right)^{-1} \simeq \frac{2(c_{L,E} - 1/2 + k)}{c_{L,E}}, \tag{13}$$

which for $c_L = 0.48$ gives $S^{-1}_{c_L} \simeq 0.16$. The remaining sensitivity above comes from the rest of the observables.

IV. FLAVOR PHYSICS

Our setup is a variation of the anarchic 5D MFV model [50,51], where the shined Yukawas are of hierarchical structure as in [36], but the SM quarks propagate in the bulk. Therefore, the following relation between the bulk masses and the 5D Yukawa matrices is obtained:

$$\begin{aligned} C_Q &= a_Q \cdot \mathbf{1}_3 + b^Q_U Y_U Y_U^\dagger + b^Q_D Y_D Y_D^\dagger + \cdots, \\ C_{U,D} &= a_{U,D} \cdot \mathbf{1}_3 + b_{U,D} Y^\dagger_{U,D} Y_{U,D} + \cdots, \end{aligned} \tag{14}$$

where the dots stand for contributions from higher powers of the Yukawa flavons. Recall also that, in the absence of mixing, the masses in terms of the Yukawas are given by

$$m_{U,D} \simeq \alpha_{U,D} \frac{v}{\sqrt{2}} F_Q Y_{U,D} F_{U,D} r^H_{00}(\beta, c_Q, c_{U,D}) + \cdots, \tag{15}$$

where F_X are matrices with eigenvalues f_{x^i} representing the IR projection of the quark zero-mode profiles, given by $f^2_{x^i} = (1 - 2c_{x^i})/(1 - \epsilon^{1-2c_{x^i}})$, c_{x^i} are the eigenvalues of the C_x matrices, $\epsilon = \exp[-\xi]$, $\xi = \log[M_{\overline{\rm Pl}}/{\rm TeV}]$, $M_{\overline{\rm Pl}}$ is the reduced Planck mass, and $r^H_{00}(\beta, c_L, c_R) \approx \frac{\sqrt{2(1+\beta)}}{2+\beta-c_L-c_R}$ is the overlap correction for a bulk Higgs [22] ($r^H_{00}(\beta, c_L, c_R) = 1$ for a brane-localized Higgs). The $\alpha_{U,D}$ coefficients are distinct from those in the expansion of Eq. (14) (in our subsequent discussion, only the combinations $\alpha_{U,D} Y_{U,D}$ appear). For simplicity we show in Eq. (15) only the part related to the zero-mode couplings and the leading term in terms of the Yukawa flavon fields. In practice, the third-generation masses are somewhat modified due to the fact that the mass eigenstates are affected by mixing with the KK-fermions, hence this is taken into account in our quantitative analysis. NDA suggests that in the most generic models $b^Q_{U,D}$, $b_{U,D}$, and $\alpha_{U,D}$ are all of order 1 in appropriate units of the curvature [51]. However, we point out that $\alpha_{U,D}$ carry different $U(1)_{Y_{U,D}, \bar{Q}, U, D, H}$ charges (which can be thought of as generalized Peccei-Quinn symmetries), and therefore a hierarchy between α_U and α_D, and between α_i and b^Q_i, b_i is natural, and can be obtained in specific models. For instance, in models of gauge-Higgs unification, α_i can be indirectly suppressed due to gauge interactions.

An immediate consequence of the MFV framework is that bounds from flavor violation in the first two generations become much weaker. This follows from an inherent suppression of right-handed currents, which require light mass insertions [39,57]. Thus, the bound from ϵ_K, which is rather strong in the anarchic case [7–11], is irrelevant here [39], since the right-handed current is suppressed by $r^4_Q m_s m_d / m^2_b$ [$r_Q \sim 2$–3; see Eq. (B4) in Appendix B] compared to the left-handed current.

As a result of the large top mass, we actually expect higher powers of the up Yukawa to be important, and these would shift the eigenvalues of the bulk masses [39,84]. The impact of top Yukawa resummation is subtle, but can be observed in flavor violation involving left-handed currents in the first two generations. This applies, in particular, to CP violation in the D system [85] (effects of order m^2_c/m^2_t are present in the kaon system, but are much harder to observe [39,86]). If the bottom Yukawa is large as well,

[8] Because of the high KK-scale observed in the semianarchic model when compared to the best fit point [see below Eq. (6)], the one d.o.f. requires us to extrapolate the results, hence the sensitivity in that case could not be computed. However, from the three d.o.f. case we see that the sensitivity is roughly the same as when the bound was compared to the $\chi^2_{\rm SM}$.

then in the presence of flavor diagonal phases, order 1 CP violating contributions are expected in $B_{d,s}$ mixing [39,87–90]. An easy way to see this is to take the two-generation limit, where the SM Lagrangian is manifestly CP conserving. In this case, higher-dimensional operators can contain a CP violating combination of the Yukawa matrices, proportional to the covariant flavor direction, $\hat{J}$ [91,92]

$$\hat{J} \propto [Y_D Y_D^\dagger, Y_U Y_U^\dagger]. \tag{16}$$

This induces CP violation in both up and down sectors, even in the two-generation case.

We shall focus on two scenarios. The first is when the bulk parameters give a small 5D bottom Yukawa. Generically, the phenomenology of this model is rather simple, and the contributions to various flavor changing processes are highly suppressed. We then slightly deform this sweet spot solution, and show how the model approaches the large bottom 5D Yukawa limit, which yields a richer flavor structure. In particular, we demonstrate how one can generate sizable new CP violating contributions in B_d and B_s mixing, and identify a natural region of the parameters where the latter dominates, as favored by the recent D0 data [93,94] and permitted by the CDF one [95] (see [40–43,96–104] for related work and [105] for a much earlier study about lepton asymmetry in the B system).

In the following we employ only a one d.o.f. analysis of the EWPT bounds, for simplicity. We also mainly focus on a comparison with the best fit point (which is now required to comply with flavor constraints).

A. Small 5D bottom Yukawa

We first analyze the flavor structure of the theory with a small bottom Yukawa. For concreteness we give an example of such a point[9],

$$\begin{aligned} C_Q &= (0.497, 0.497, 0.348),\\ C_U &= (-0.5, -0.5, 0.482),\\ C_D &= (0.56, 0.56, 0.55),\\ \alpha_U Y_U &= (3.6\times 10^{-5}, 0.017, 6.2),\\ \alpha_D Y_D &= (0.0013, 0.024, 0.36). \end{aligned} \tag{17}$$

The reader should bear in mind though, that as long as the bottom Yukawa is small, the gross features of the model near the sweet spot remain unchanged. The resulting bound from EWPTs is 2.4 TeV.

In the limit of small Y_D, the bulk masses can be expanded in powers of Y_U only. This is manifest in the choice of bulk masses in Eq. (17), where C_D is almost completely diagonal, and in $C_{Q,U}$ only the third eigenvalue is shifted away from the first two. Since this applies to the F_X's as well, we have according to Eq. (15) $[m_U, Y_U] = 0$, i.e., m_U and Y_U can be simultaneously diagonalized. Our model thus contains a built-in up-type flavor alignment, hence FCNCs are only present in the down sector. Moreover, flavor violation in the down sector is proportional to elements of the Cabibbo-Kobayashi-Maskawa (CKM) matrix $V^{\rm CKM}$, and right-handed currents are significantly suppressed, as anticipated since our current setup belongs to the MFV framework with a small bottom Yukawa.

It is crucial to emphasize that when expanding the bulk masses as functions of the Yukawa matrices, higher order terms in Y_U are important, and may give rise to a significant effect, as shown below. Therefore, we write

$$\begin{aligned} C_Q &= a_Q\cdot \mathbf{1}_3 + b_U^Q Y_U Y_U^\dagger + b_D^Q Y_D Y_D^\dagger + d_{UU}^Q Y_U (Y_U Y_U^\dagger) Y_U^\dagger \\ &\quad + d_{DU}^Q Y_D (Y_U Y_U^\dagger) Y_D^\dagger + \cdots,\\ C_U &= a_U\cdot \mathbf{1}_3 + b_U Y_U^\dagger Y_U + d_{UU} Y_U^\dagger (Y_U Y_U^\dagger) Y_U + \cdots,\\ C_D &= a_D\cdot \mathbf{1}_3 + b_D Y_D^\dagger Y_D + d_{DU} Y_D^\dagger (Y_U Y_U^\dagger) Y_D + \cdots, \end{aligned} \tag{18}$$

where some of these terms are actually small in the case of small bottom Yukawa.

The most severe constraints are from the $B_{d,s}$ systems, in the form of a $\Delta B = 2$ contribution to the mixing amplitude. These are generated in RS via a tree-level KK-gluon exchange, formulated in terms of two of the standard four-quark operators,

$$Q_1 = \bar{q}_{jL}^\alpha \gamma_\mu q_{iL}^\alpha \bar{q}_{jL}^\beta \gamma_\mu q_{iL}^\beta, \quad Q_4 = \bar{q}_{jR}^\alpha q_{iL}^\alpha \bar{q}_{jL}^\beta q_{iR}^\beta, \tag{19}$$

where α, β are color indices and i, j are flavor indices. New physics in the $B_{d,s}$ mixing amplitudes can be described by four real parameters,

$$M_{12}^{d,s} = (M_{12}^{d,s})^{\rm SM}(1 + h_{d,s} e^{2i\sigma_{d,s}}), \tag{20}$$

where M_{12} is the dispersive part of the amplitude. We shall use the notation $h_{d,s}^{1,4}$, where the superscript denotes the contributing operator.

In order to evaluate the flavor-violating contribution to B_d, we need to rotate the diagonal coupling of two quarks with the KK-gluon to the mass basis. Since the mass basis is aligned with Y_U, this introduces CKM factors (plus subleading corrections for large bottom Yukawa) in the case of left-handed quarks, and a factor related to the difference of overlaps of the b and d quarks with the KK-gluon. This calculation is performed in detail in Appendix B [see Eq. (B10)]. The Wilson coefficient for Q_1 is then given by

$$C_1 \approx \frac{g_{s*}^2}{6m_{\rm KK}^2}(V_{tb}V_{td}^*)^2[f_{Q^3}^2 r_{00}^g(c_{Q^3}) - f_{Q^1}^2 r_{00}^g(c_{Q^1})]^2. \tag{21}$$

Here g_{s*} is the dimensionless 5D coupling of the gluon ($g_{s*} = 3$ with one-loop matching), $r_{00}^g(c) \approx \frac{\sqrt{2}}{J_1(x_1)} \frac{0.7}{6-4c} \times (1 + e^{c/2})$ is the overlap correction for two zero-mode

[9]In the context of flavor physics, it is more convenient to employ the notations c_{U^3} and c_{D^3} instead of c_t and c_b used above. We thus adopt this change in this section.

quarks with the KK-gluon [9,20,22], with $x_1 \cong 2.4$ as the first root of the Bessel function $J_0(x_1) = 0$, and $(V_{tb}V_{td}^*)^2 \approx [V_{tb}^{\rm CKM}(V_{td}^{\rm CKM})^*]^2(1 + rY_b^2 e^{i2\theta_d})$, with θ_d an arbitrary phase and r a proportionality coefficient (in the current case we neglect this correction, which is formally of order Y_b^2). A similar formula applies for B_s (replacing $d \to s$ and $1 \to 2$).

For a right-handed coupling, which is a part of the Q_4 contribution, the rotation is more involved, and introduces some additional factors (see Appendix B). The resulting Wilson coefficient is

$$C_4 \approx \frac{g_{s*}^2}{m_{\rm KK}^2}(V_{tb}V_{td}^*)^2 \frac{m_d}{m_b}\left[\left(\frac{f_{Q^3} r_{00}^H(\beta, c_{Q^3}, c_{D^3})}{f_{Q^1} r_{00}^H(\beta, c_{Q^1}, c_{D^3})}\right)^2 - 1\right] \times [f_{Q^3}^2 r_{00}^g(c_{Q^3}) - f_{Q^1}^2 r_{00}^g(c_{Q^1})][f_{D^3}^2 r_{00}^g(c_{D^3}) - f_{D^1}^2 r_{00}^g(c_{D^1})]. \tag{22}$$

In order to derive a bound on the KK-scale, we allow for $h_d^{1,4}$ to be as large as 0.5 (since the NP contributions do not carry additional CP phases) [40]. We include running and mixing effects at 2 TeV, as described in [7] and references therein. The bound resulting from Q_1 is

$$\left(\frac{m_{\rm KK}}{2\ \rm TeV}\right) \gtrsim 3.7(\delta f_{Q^{31}}^2) \approx 2.3\left(\frac{1 - 2.1c_{Q^3}}{1 - \frac{2}{3}c_{Q^3}}\right), \tag{23}$$

where we defined

$$(\delta f_{Q^{ij}}^2) \equiv f_{Q^i}^2 r_{00}^g(c_{Q^i}) - f_{Q^j}^2 r_{00}^g(c_{Q^j}), \tag{24}$$

and used $c_{Q^1} = 0.497$ from Eq. (17) and

$$f_x^2 r_{00}^g(c_x) \approx \frac{1 - 2c_x}{1.5 - c_x}, \tag{25}$$

which is a good approximation for $0 < c_x < 0.47$. Note that for $c_{Q^3} = 0.348$ we have $m_{\rm KK} \gtrsim 1.9$ TeV, consistent with EWPT.[10] Similarly, the bound from Q_4 is

$$\left(\frac{m_{\rm KK}}{2\rm TeV}\right) \gtrsim 23\sqrt{\frac{m_d}{m_b}\left[\left(\frac{f_{Q^3} r_{00}^H(\beta, c_{Q^3}, c_{D^3})}{f_{Q^1} r_{00}^H(\beta, c_{Q^1}, c_{D^3})}\right)^2 - 1\right](\delta f_{Q^{31}}^2)(\delta f_{D^{31}}^2)}. \tag{26}$$

The actual constraint is much weaker than Eq. (23) because of the m_d/m_b suppression and the approximate degeneracy of the f_D's. It is instructive to see the relation between the contributions of Q_4 and Q_1 to B_d mixing,

$$\left.\frac{C_4}{C_1}\right|_{2\ \rm TeV} \approx 40\frac{m_d}{m_b}\frac{(\delta f_{D^{31}}^2)}{(\delta f_{Q^{31}}^2)}\left[\left(\frac{f_{Q^3} r_{00}^H(\beta, c_{Q^3}, c_{D^3})}{f_{Q^1} r_{00}^H(\beta, c_{Q^1}, c_{D^3})}\right)^2 - 1\right]. \tag{27}$$

The same exercise can be carried out for B_s mixing, where now we allow the RS contribution to be 30% of the SM one (that is, $h_s^{1,4} = 0.3$), without new phases [40]. The bounds from Q_1 and Q_4 are

$$\left(\frac{m_{\rm KK}}{2\ \rm TeV}\right) \gtrsim 4.7(\delta f_{Q^{32}}^2) \approx 3\left(\frac{1 - 2.1c_{Q^3}}{1 - \frac{2}{3}c_{Q^3}}\right),$$
$$\left(\frac{m_{\rm KK}}{2\ \rm TeV}\right) \gtrsim 30\sqrt{\frac{(\delta f_{Q^{31}}^2)}{f_{Q^1}^2 r_{00}^g(c_{Q^1})}(\delta f_{Q^{32}}^2)(\delta f_{D^{32}}^2)\frac{m_s}{m_b}}, \tag{28}$$

respectively. For $c_{Q^3} = 0.348$ the first bound reads $m_{\rm KK} \gtrsim 2.4$ TeV. The Q_4 bound is much stronger than for B_d, but still weaker than the one from Q_1. Note that the Q_1 contribution is universal, i.e., the same for B_d and B_s, and that the bound in the first line of Eq. (28) is stronger than Eq. (23) only because we required $h_d = 0.5$ and $h_s = 0.3$. Equation (27) changes for B_s to

$$\left.\frac{C_4}{C_1}\right|_{2\ \rm TeV} \approx 39\frac{m_s}{m_b}\frac{(\delta f_{D^{32}}^2)}{(\delta f_{Q^{32}}^2)}\left[\left(\frac{f_{Q^3} r_{00}^H(\beta, c_{Q^3}, c_{D^3})}{f_{Q^2} r_{00}^H(\beta, c_{Q^2}, c_{D^3})}\right)^2 - 1\right]. \tag{29}$$

Note that in our example $(\delta f_{Q^{31}}^2) = (\delta f_{Q^{32}}^2)$.

One may wonder whether $\Delta B = 1$ processes, such as $b \to s\gamma$, could also lead to significant bounds on the model (see e.g. [18,22] for some recent estimations within the anarchic scenario). However, since this is a chirality-flipping process, it must involve right-handed mixing angles, which are strongly suppressed in our model, as shown in Eq. (B9). More generally, this is a consequence of the fact that our model belongs to the class of general MFV [39], where right-handed currents are suppressed by a ratio of masses, that is m_s/m_b in this case.

To summarize this example, characterized by Eq. (17), the overall bound that we find is

$$m_{\rm KK} \gtrsim 2.4\ \rm TeV, \tag{30}$$

coming from the Q_1 contribution to B_s and from EWPTs. It should be noted that this bound can be reduced to 2.2 TeV if compared to the SM, instead of the best fit point (with an appropriate change in the bulk masses).

[10] This is weaker than in [106], which was ultraconservative.

B. Large 5D bottom Yukawa

The analysis of the previous subsection assumed a small bottom Yukawa, as can be inferred from Eq. (17). Yet by reducing α_D, for example, the bottom Yukawa can be made larger, until it is of order 1. Consequently, Y_D resummation effects appear, and the results of the previous subsection receive $\mathcal{O}(1)$ corrections plus a general phase [39].

We can try to use the large bottom Yukawa case to obtain a larger RS contribution to B_s than for B_d. Since C_1 is universal in that sense, this requires one to increase C_4 to be larger than C_1, noting that $h_s^4 > h_d^4$ in any case.

Considering as an example the following bulk masses,

$$\begin{aligned} C_Q &= (0.516, 0.516, 0.35), \\ C_U &= (-0.5, -0.5, 0.479), \\ C_D &= (0.56, 0.56, 0.497), \\ \alpha_U Y_U &= (5.1 \times 10^{-5}, 0.025, 5.9), \\ \alpha_D Y_D &= (0.0018, 0.034, 0.12), \end{aligned} \tag{31}$$

and an appropriate α_D to obtain a large bottom Yukawa, we have the following results:

(i) The bound on the KK-scale from EWPT is slightly raised to 2.6 TeV.

(ii) Because of the generic phase, it is required to take $h_d^{1,4}$ to be 0.3 instead of 0.5 [40].

(iii) As a result of taking $c_{D^3} = 0.496$, we now have $h_s^4 \cong 1.33 h_s^1 \cong 0.4$, while for the B_d system C_4 is still smaller than C_1 [see Eqs. (27) and (29), when evaluated at the scale 2.6 TeV originating from EWPT constraints].

(iv) Another possible point is $c_{D^3} = 0.4$ (and some more slight adjustments of other bulk masses). Then the EWPT bound is ~2.7 TeV and $h_s^4 \cong 1.75$.

The implication of this result is that our model is now in accordance with the recent Tevatron data, which favor larger contributions to B_s than for B_d [40]. The price to pay is that $\alpha_D Y_b \approx 0.12$, so that in order to have an $\mathcal{O}(1)$ bottom Yukawa, α_D must be small. While this is technically natural, it still requires a small parameter to be tuned by hand.

It is actually simple to explain why our model cannot produce $h_s > 0.3$ and $h_d \lesssim 0.3$ if we insist on having a large bottom Yukawa with $\alpha_D = \mathcal{O}(1)$. The latter requirement leads to the relation $f_{Q^3} f_{D^3} \lesssim 0.01$, in order to get the correct bottom mass. However, as can be seen from Eq. (22), the C_4 contribution is roughly proportional to $(f_{Q^3} f_{D^3})^2$ (times another factor of $f_{Q^3}^2$ which is smaller than 1), and as a result it is too small to yield $h_s > 0.3$.

The universal $h_d = h_s$ *case*

While the data favor large *CP* violation in the B_s system, a reasonable fit of the flavor measurements is obtained in the $SU(2)$ universal case where $h_b \equiv h_d = h_s \sim 0.3$, consistent with the data [40]. It is not surprising that our framework (as well as the anarchic RS case [4,5]) can account for this case in a straightforward manner, while having α_D and the 5D bottom Yukawa of order unity. This is obtained by taking c_{D^3} to be ~0.6 and c_{D^i} around 0.6–0.65, while the other bulk masses are as in Eq. (17). In this case, one can sharply predict order 1 CPV phases with exact universality, $\sigma_b \equiv \sigma_d = \sigma_s$ [39]. The resulting EWPT bound is ~2.4 TeV.

C. Higgs mediated FCNCs

Another possible source of flavor violation arises from the Higgs [107–109], which obtains off-diagonal couplings in the mass basis as a result of mixing between zero mode and KK-fermions. For an IR brane Higgs, the leading spurion which induces this process is [108]

$$\sim F_Q Y_D Y_D^\dagger Y_D F_D^\dagger, \tag{32}$$

omitting all universal factors.[11] Yet, the resulting flavor violation is suppressed relative to the KK-gluon contribution. To see this, let us neglect the masses (and Yukawa couplings) of quarks of the first two generations. Then in its diagonal basis, Y_D is proportional to diag$(0, 0, Y_b)$, and consequently we have $Y_D^3 \propto Y_b^2 Y_D$. In other words, the leading mass term in Eq. (15) and the spurion in Eq. (32) are aligned together, so no flavor violation is generated. Restoring the strange mass, we expect to have a $(m_s/m_b)^2$ suppression, after squaring these spurions to obtain the relevant Wilson coefficients. Since a factor of this kind does not appear for the KK-gluon contribution to flavor violation via Q_1, the Higgs effect can be neglected.

This argument is easily generalized to the bulk Higgs case. The Y_D^3 part of Eq. (32) should be written now as

$$Y_D r_{01}^H Y_D^\dagger r_{10}^H Y_D, \tag{33}$$

where $r_{01,10}^H$ is an overlap correction for the coupling of the Higgs to a zero-mode quark and a KK-quark. Even though these corrections are not universal, the wrapping Y_D's act as a projection operator for the 3-3 matrix element, when neglecting the first two generations' masses. Therefore, we still have $Y_D^3 \propto Y_b^2 Y_D$, and the conclusion from before applies to this case as well. Moreover, we did not have to assume anything about F_Q and F_D, hence the Higgs contribution is negligible in both the small and the large bottom Yukawa cases.

V. CONCLUSIONS

We analyzed a warped 5D model where the SM Yukawa hierarchy is set by UV physics, which realizes a bulk

[11]An additional contribution comes from a one-loop process involving a charged Higgs and up-type quarks. However, as a result of the loop suppression, it is subleading.

TABLE II. Bounds (in TeV) from EWPTs for the various statistical scenarios considered, comparing the flavor triviality model to the semianarchic case. Best fit refers to the bound relative to the best fit point, where χ^2 is lower than in the SM, and SM refers to the case where the SM is assumed to be the minimum χ^2. In the last two columns the flavor-triviality analysis is for six d.o.f., while the semianarchic one is for three d.o.f.

	One d.o.f.		Six/three d.o.f.	
Model	Best fit	SM	Best fit	SM
Flavor triviality	2.1	1.8	1.7	1.6
Semianarchy	7	4.6	4.6	3.9

version of the Rattazzi-Zaffaroni model [36]. Such a scenario displays the weakest bound on the scale of RS type of new physics explored to date, both from the point of view of electroweak precision measurements, as well as from flavor constraints. The EW precision tests allow for a sweet spot with a KK-scale as low as 2.1 TeV, which is more than a factor of 2 lower than in the anarchic RS setup [with fermions in the minimal $SU(2)_L \times SU(2)_R$ representations, as discussed in the main text]. Such a low scale for the RS KK physics should lead to significantly better prospects for discovery at the LHC, given the fact that its reach for a KK-gluon is around 4 TeV [110,111]. A summary of the bounds that we find from EWPTs for different statistical scenarios is presented in Table II.

This model, by construction, belongs to the minimal flavor violation (MFV) framework [52–61], and naively one would expect a rather dull flavor phenomenology. Indeed the model is strongly protected from CP violation in the first two generations. Imposing the flavor constraints, we also find consistency with a KK-scale of about 2.4 TeV. Thus the RS ϵ_K problem is avoided, practically, without interfering with the model's visibility. This is a natural consequence of MFV. However, the fact that in this framework the third-generation couplings are sizable and flavor-violating couplings effectively exponentiate leads to various interesting deviations from the commonly studied MFV models. Thus, this class of models belongs to the *general* MFV framework [39]. Performing a deformation around the best fit point in parameter space allows for a rather rich third-generation flavor phenomenology, such as providing the new CPV source required by the latest same-sign dimuon signal from D0. The present ideas are expected to help reduce the constraints (hence the fine-tuning) of extra-dimensional scenarios of EWSB, such as models where the Higgs is a pseudo-Goldstone boson. The detailed study of such an exciting possibility is left to future work.

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APPENDIX A: MATCHING RS TO THE EW PRECISION OPERATORS

New physics effects at the weak scale are captured by a set of effective operators added to the renormalizable part of the SM Lagrangian: $\mathcal{L} = \mathcal{L}_{\rm SM} + \sum_i a_i \mathcal{O}_i$, where $\mathcal{O}_i$ are gauge and flavor invariant operators. In the absence of flavor and CP violation, 20 operators[12] (of mass dimension-6) contribute most significantly to the electroweak precision observables [68]. There are 2 operators affecting the gauge sector,

$$\mathcal{O}_{WB} = h^\dagger \sigma^a h W^a_{\mu\nu} B^{\mu\nu}, \quad \mathcal{O}_h = |h^\dagger D_\mu h|^2, \tag{A1}$$

which generate, respectively, the S and T parameters, 7 operators shifting the fermion-gauge boson couplings

$$\begin{aligned} \mathcal{O}^s_{hf} &= i h^\dagger D_\mu h \bar{f} \gamma^\mu f + \text{H.c.}, \\ \mathcal{O}^s_{hF} &= i h^\dagger D_\mu h \bar{F} \gamma^\mu F + \text{H.c.}, \\ \mathcal{O}^t_{hF} &= i h^\dagger \sigma^a D_\mu h \bar{F} \gamma^\mu \sigma^a F + \text{H.c.}, \end{aligned} \tag{A2}$$

where $f = u$, d, e and $F = q$, l, and 11 four-fermion operators contributing to the leptonic sector

$$\begin{aligned} \mathcal{O}^s_{ll} &= \frac{1}{2}(\bar{l}\gamma^\mu l)^2, & \mathcal{O}^t_{ll} &= \frac{1}{2}(\bar{l}\gamma^\mu \sigma^a l)^2, \\ \mathcal{O}^s_{le} &= (\bar{l}\gamma^\mu l)(\bar{e}\gamma_\mu e), & \mathcal{O}^s_{ee} &= \frac{1}{2}(\bar{e}\gamma^\mu e)^2, \\ \mathcal{O}^s_{lq} &= (\bar{l}\gamma^\mu l)(\bar{q}\gamma_\mu q), & \mathcal{O}^t_{lq} &= (\bar{l}\gamma^\mu \sigma^a l)(\bar{q}\gamma_\mu \sigma^a q), \\ \mathcal{O}^s_{qe} &= (\bar{q}\gamma^\mu q)(\bar{e}\gamma_\mu e), & \mathcal{O}^s_{lu} &= (\bar{l}\gamma^\mu l)(\bar{u}\gamma_\mu u), \\ \mathcal{O}^s_{ld} &= (\bar{l}\gamma^\mu l)(\bar{d}\gamma_\mu d), & \mathcal{O}^s_{eu} &= (\bar{e}\gamma^\mu e)(\bar{u}\gamma_\mu u), \end{aligned}$$
$$\mathcal{O}^s_{ed} = (\bar{e}\gamma^\mu e)(\bar{d}\gamma_\mu d). \tag{A3}$$

Whenever relevant, a $U(3)$ trace over flavor is assumed in all of the above. Given the peculiar behavior of the third-generation quarks, new physics is expected to break the $U(3)^3$ flavor symmetries in the quark sector down to $[U(2) \times U(1)]^3$. In our setup b_R behaves as the lighter generations d_R and s_R, and EWPTs are not sensitive to top observables. Thus, it is well justified to work in a limit where only $U(3)_Q$ is broken down to $U(2)_q \times U(1)_Q$, with

[12] An additional operator, $\mathcal{O}_W = \epsilon_{abc} W^a_{\mu\nu} W^{\nu\rho b} W^{\mu c}_\rho$, can be considered as well. However, it is weakly constrained by EWPT, since it affects only the triple and quadruple gauge self-couplings, which are poorly measured. Thus we set this operator to zero in our fit.

q and Q denoting the first two- and third-generation quark doublets, respectively. In this case there are 5 additional operators

$$\begin{aligned}
\mathcal{O}^s_{hQ} &= ih^\dagger D_\mu h \bar{Q}\gamma^\mu Q + \text{H.c.},\\
\mathcal{O}^t_{hQ} &= ih^\dagger \sigma^a D_\mu h \bar{Q}\gamma^\mu \sigma^a Q + \text{H.c.},\\
\mathcal{O}^t_{lQ} &= (\bar{l}\gamma^\mu \sigma^a l)(\bar{Q}\gamma_\mu \sigma^a Q),\\
\mathcal{O}^s_{lQ} &= (\bar{l}\gamma^\mu l)(\bar{Q}\gamma_\mu Q),\\
\mathcal{O}^s_{Qe} &= (\bar{Q}\gamma^\mu Q)(\bar{e}\gamma_\mu e),
\end{aligned} \tag{A4}$$

and a $U(2)$ trace over the first two generations is now understood in the $\bar{q}\gamma^\mu q$ current of $\mathcal{O}^{s,t}_{hq}$, $\mathcal{O}^{s,t}_{lq}$, and $\mathcal{O}^s_{qe}$. Therefore, 25 operators are relevant for EWPT. We use the SM global fit of [46], following the approach developed in [68,69], with updated $m_{\text{top}} = (173.3 \pm 1.1)$ GeV [112] and $m_W = (80.420 \pm 0.031)$ GeV [113] measurements from the Tevatron.

a. Tree-level effects.—We start with matching the coefficients of the 25 operators to RS at tree-level. The leading contributions arise from exchange of gauge KK-modes, as depicted in the diagrams of Fig. 1. An explicit evaluation of the latter yields (see [65] for a pedagogical review[13])

$$\begin{aligned}
a_h &= \frac{g_5'^2}{2}(G_{++} - G_{-+}),\\
a^t_{hF} &= \frac{g_5^2}{4} I_{++}(c_F),\\
a^t_{FF'} &= \frac{g_5^2}{4} J_{++}(c_F, c_{F'}),\\
a^s_{hF} &= \frac{g_5'^2}{2} Y_F[I_{++}(c_F) - I_{-+}(c_F)] + \frac{g_{5R}^2}{2} T^F_{3R} I_{-+}(c_F),\\
a^s_{hf} &= \frac{g_5'^2}{2} Y_f[I_{++}(c_f) - I_{-+}(c_f)] + \frac{g_{5R}^2}{2} T^f_{3R} I_{-+}(c_f),\\
a^s_{FF'} &= g_5'^2 Y_F Y_{F'} J_{++}(c_F, c_{F'}) + \frac{g_{5R}^2}{\cos^2\theta}(T^F_{3R} - \sin^2\theta Y_F)\\
&\quad \times (T^{F'}_{3R} - \sin^2\theta Y_{F'}) J_{-+}(c_F, c_{F'}),\\
a^s_{ff'} &= g_5'^2 Y_f Y_{f'} J_{++}(c_f, c_{f'}) + \frac{g_{5R}^2}{\cos^2\theta}(T^f_{3R} - \sin^2\theta Y_f)\\
&\quad \times (T^{f'}_{3R} - \sin^2\theta Y_{f'}) J_{-+}(c_f, c_{f'}),
\end{aligned} \tag{A5}$$

where $F, F' = Q, q, l$ and $f, f' = u, d, e$ with $\sin^2\theta = g_5'^2/g_{5R}^2$. The G, I, J wave-function overlap integrals are given by

$$\begin{aligned}
G_{\pm+} &= v^{-4}\int_R^{R'} dz dz' \left(\frac{R}{z}\right)^3\\
&\quad \times \left(\frac{R}{z'}\right)^3 v_5(z,\beta)^2 \mathcal{G}_{\pm+}(z,z') v_5(z',\beta)^2,\\
I_{\pm+}(c) &= v^{-2}\int_R^{R'} dz dz' \left(\frac{R}{z}\right)^4\\
&\quad \times \left(\frac{R}{z'}\right)^3 \chi(z,c)^2 \mathcal{G}_{\pm+}(z,z') v_5(z',\beta)^2,\\
J_{\pm+}(c,c') &= \int_R^{R'} dz dz' \left(\frac{R}{z}\right)^4\\
&\quad \times \left(\frac{R}{z'}\right)^4 \chi(z,c)^2 \mathcal{G}_{\pm+}(z,z') \chi(z',c')^2,
\end{aligned} \tag{A6}$$

where $\mathcal{G}_{\pm+}$ is the mixed position-momentum 5D propagator for $(\pm, +)$ gauge bosons in anti–de Sitter space evaluated at zero (4D) momentum [74,114] and $\chi(z, c)$ is the fermion zero-mode wave function, while $v_5(z, \beta)$ is the bulk Higgs VEV. g_5, g_5', and g_{5R} are the 5D gauge couplings of $SU(2)_L$, $U(1)_Y$, and $SU(2)_R$, respectively. While g_5 and g_5' have to be matched to the 4D gauge couplings (see below), g_{5R} is a free parameter of the model, which we take to be $g_{5R} = g_5$, as required by the extended custodial symmetry that protects the $Zb_L\bar{b}_L$ coupling. Note that $a_{WB} \sim \mathcal{O}(v^4/m^4_{\text{KK}})$ at tree level, and we recall that the tree-level S parameter often quoted in RS is coming from a "universal" shift to the fermion couplings. This contribution is included in the global fit through the shifts in the fermion-gauge boson couplings, which is just a consequence of the fact that some operators in the effective Lagrangian are redundant [115,116].

b. Matching of 5D gauge couplings.—The 5D gauge couplings used above have to be matched to their 4D values in the effective action. Including one-loop renormalization, the matching conditions are [74–76,117–123]

$$\frac{1}{g^2} = \log(R'/R)\left(\frac{1}{g_5^2 k} + \frac{b_g}{8\pi^2}\right) + \frac{1}{g^2_{\text{UV}}} + \frac{1}{g^2_{\text{IR}}}, \tag{A7}$$

where the last two terms are contributions from (possible) "bare" brane-localized kinetic terms, which we set to zero for simplicity. The one-loop β-function coefficients b receive contributions from the bulk only through elementary fields. Hence after removing the Higgs contribution from the running, we find $b_g = -10/3$ and $b_{g'} = 20/3$. Therefore, matching the 5D gauge couplings at the TeV scale yields $g_5\sqrt{k} \simeq 27.3 g/\sqrt{\log(R'/R)} \simeq 3.02$ and $g_5'\sqrt{k} = 43.9 g'/\sqrt{\log(R'/R)} \simeq 2.66$ for $k = R^{-1} = 2.4 \times 10^{18}$ GeV and $m_{\text{KK}} \sim 2$ TeV.

c. One-loop effects.—The large top Yukawa induces a non-negligible contribution to the S and T parameters, as well as to the b_L coupling to Z. A straightforward calculation of the one-loop diagram of Fig. 2 gives the following contributions to the oblique parameters [27,124]:

[13]For reference, the relation between the notation used here and that of [65] is as follows: $\alpha^N = LG_{++}$, $\alpha^D = LG_{-+}$, $\beta^N_\psi = LI_{++}(c_\psi)$, $\beta^D_\psi = LI_{-+}(c_\psi)$, $\gamma^N_{\psi\psi'} = LJ_{++}(c_\psi, c_{\psi'})$, and $\gamma^D_{\psi\psi'} = LJ_{-+}(c_\psi, c_{\psi'})$, where $L = R\log[R'/R]$ is the proper size of the fifth dimension.

$$S=\frac{N_c}{2\pi}\sum_{\alpha,\beta}\sum_{X=U,D}[(X^{L\dagger}_{\alpha\beta}Y^{L}_{X\beta\alpha}+X^{R\dagger}_{\alpha\beta}Y^{R}_{X\beta\alpha})\chi_+(m^\alpha_X,m^\beta_X)$$
$$+(X^{L\dagger}_{\alpha\beta}Y^{R}_{X\beta\alpha}+X^{R\dagger}_{\alpha\beta}Y^{L}_{X\beta\alpha})\chi_-(m^\alpha_X,m^\beta_X)],$$
$$T=\frac{N_c}{16\pi s_W^2c_W^2m_Z^2}\bigg[\sum_{\alpha,i}V^2_{\alpha i}\theta_+(m^\alpha_U,m^i_D)+A^V_{\alpha i}\theta_-(m^\alpha_U,m^i_D)$$
$$-\sum_{\beta<\alpha}U^2_{\alpha\beta}\theta_+(m^\alpha_U,m^\beta_U)+A^U_{\alpha\beta}\theta_-(m^\alpha_U,m^\beta_U)\bigg], \tag{A8}$$

where we defined $K^2\equiv|K^L|^2+|K^R|^2$, $A^K=2\,\mathrm{Re}[K^LK^{R*}]$ with $K=U$, V. The unitary matrices $U^{L,R}$,$D^{L,R}$ ($Y^{L,R}_{U,D}$) denote the couplings of the $Q=2/3$ and $Q=(-1/3,5/3)$ mass eigenstates to the W^μ_{3L} (B^μ) zero mode, while the $V^{L,R}$ matrices stand for the coupling of the mass eigenstates to the $W^\pm$ zero mode. The definitions of the loop functions $\theta_\pm$ and $\chi_\pm$ are [124]

$$\theta_+(m_1,m_2)=m_1^2+m_2^2-\frac{2m_1^2m_2^2}{m_1^2-m_2^2}\log\frac{m_1^2}{m_2^2}, \tag{A9}$$

$$\theta_-(m_1,m_2)=2m_1m_2\left(\frac{m_1^2+m_2^2}{m_1^2-m_2^2}\log\frac{m_1^2}{m_2^2}-2\right), \tag{A10}$$

$$\chi_+(m_1,m_2)=\frac{5(m_1^4+m_2^4)-22m_1^2m_2^2}{9(m_1^2-m_2^2)^2}-\frac{2}{3}\log\frac{m_1m_2}{\mu^2}$$
$$+\frac{3m_1^2m_2^2(m_1^2+m_2^2)-(m_1^6+m_2^6)}{3(m_1^2-m_2^2)^3}\log\frac{m_1^2}{m_2^2}, \tag{A11}$$

$$\chi_-(m_1,m_2)=\frac{m_1m_2}{(m_1^2-m_2^2)^3}\left(m_1^4-m_2^4-2m_1^2m_2^2\log\frac{m_1^2}{m_2^2}\right), \tag{A12}$$

where the renormalization scale dependence in χ_+ cancels out in S thanks to $\mathrm{tr}[U^\dagger Y_U+D^\dagger Y_D]=0$. Note that Eq. (A8) includes SM contributions from top and bottom

$$S_{\rm SM}\simeq\frac{N_c}{18\pi}\left[3-\log\left(\frac{m_t^2}{m_b^2}\right)\right],\qquad T_{\rm SM}\simeq\frac{N_c}{16\pi s_W^2c_W^2}\left(\frac{m_t^2}{m_Z^2}\right), \tag{A13}$$

which need to be subtracted in order to isolate the new physics contributions.

These loop effects are controlled by EWSB, dominantly from the top sector as parametrized by the 5D top Yukawa. The contributions associated with EWSB mixing among heavy KK-modes are controlled by the top Yukawa coupling evaluated at a scale of the order of the KK masses. These contributions are subdominant, however, and the result is dominated by EW mixing between the KK-modes and the top zero-mode. The relevant diagrams display an IR divergence that is cut off by the top mass, indicating that the result is dominated by scales of order $\mu\sim m_{\rm top}$ (see [27]). To be conservative, we use in these loop contributions a top running mass $m_t(\mu=m_{\rm top})\approx 160$ GeV [125]. Similarly, we use $m_b(\mu=m_{\rm top})\approx 2.7$ GeV.

In gauge-Higgs unified models, where the Higgs is realized as a pseudo-Goldstone boson, the Higgs only partially regulates the divergent contribution to the S and T parameters arising from loops of (longitudinal) SM gauge fields[14] [70]. This results in a logarithmic correction to the S and T parameters which is cut by the KK-scale [71],

$$\Delta S=\frac{1}{12\pi}(1-a^2)\log\frac{\Lambda^2_{\rm eff}}{m_h^2},\qquad \Delta T=-\frac{3}{16\pi c_W^2}(1-a^2)\log\frac{\Lambda^2_{\rm eff}}{m_h^2}, \tag{A14}$$

where $\Lambda_{\rm eff}\simeq m_{\rm KK}$ and a measures the amount of Higgs compositeness, with $a=1$ corresponding to the fully elementary SM Higgs. Deviation from $a=1$ also leads to an incomplete unitarization of the W/Z scattering amplitude, and perturbative unitarity is lost at $\Lambda_{\rm eff}\simeq 1.2\ {\rm TeV}/\sqrt{|1-a^2|}$. Requiring unitarity not to be violated below the KK-scale, we find the contributions in (A14) to raise the bound for one (six) d.o.f. by 300 (200) GeV, assuming the SM as the best fit.

The $Z\bar{b}_Lb_L$ vertex also receives large radiative corrections dominated by the diagrams of Fig. 3. This yields

$$\delta g^{b_L}=\frac{\alpha}{2\pi}\bigg[\sum_\alpha V^L_{\alpha b}V^L_{\alpha b}\bigg[F_{\rm SM}(r_\alpha)+\tilde F\left(\frac{U^L_{\alpha\alpha}}{2}-\frac{1}{2},\frac{U^R_{\alpha\alpha}}{2},r_\alpha\right)\bigg]+\sum_{\alpha<\beta}V^L_{\alpha b}V^L_{\beta b}\mathcal{F}\left(\frac{U^L_{\alpha\beta}}{2},\frac{U^R_{\alpha\beta}}{2},r_\alpha,r_\beta\right)\bigg], \tag{A15}$$

where $r_\alpha=(m^\alpha_U/m_W)^2$ and the loop functions are [45,126]

$$F_{\rm SM}(r)=\frac{r}{8s_W^2}\frac{(r-1)(r-6)+(3r+2)\log r}{(r-1)^2}, \tag{A16}$$

$$\tilde F(g_L,g_R,r)=\frac{r}{8s_W^2}\bigg[g_L\left(2-\frac{4}{r-1}\log r\right)-g_R\left(\frac{2r-5}{r-1}+\frac{r^2-2r+4}{(r-1)^2}\log r\right)\bigg], \tag{A17}$$

[14] We thank Kaustubh Agashe for bringing this point to our attention.

$$\mathcal{F}(g_L, g_R, r, r') = \frac{1}{4s_W^2(r'-r)}\left[2g_L\left(\frac{r-1}{r'-1}r'^2\log r' - \frac{r'-1}{r-1}r^2\log r\right) - g_R\sqrt{rr'}\left(r'-r+\frac{r'-4}{r'-1}r'\log r' - \frac{r-4}{r-1}r\log r\right)\right]. \quad \text{(A18)}$$

Here again the SM contribution,

$$\delta g_{\rm SM}^{b_L} = \frac{\alpha}{2\pi}F_{\rm SM}(r_t), \quad \text{(A19)}$$

should be subtracted to isolate the contribution from new physics. Note that this result is derived for an off-shell ($q^2 = 0$) Z in the 't Hooft/Feynman ($\xi = 1$) gauge. Although the result should only be gauge-invariant when the Z is on-shell ($q^2 = m_Z^2$), we expect the missing terms to suffer an additional $m_Z^2/m_{\rm KK}^2$ suppression, so the $q^2 = 0$ result quoted constitutes a valid approximation for the new physics contribution. Notice also that all the radiative corrections above decouple like $v^2/m_{\rm KK}^2$, as expected, since they arise from vectorlike (KK-)fermions which mix with the chiral zero-mode through Yukawa couplings.

The above one-loop corrections are accounted for in the global fit by adding the following shifts to the coefficients of the $\mathcal{O}_{WB}$, $\mathcal{O}_h$, and $\mathcal{O}_{hQ}^s$ operators:

$$\begin{aligned} a_{WB} &\to a_{WB} + \frac{gg'}{16\pi v^2}(S - S_{\rm SM}),\\ a_h &\to a_h - \frac{g^2 s_W^2}{2\pi v^2}(T - T_{\rm SM}),\\ a_{hQ}^s &\to a_{hQ}^s - \frac{2}{v^2}(\delta g^{b_L} - \delta g_{\rm SM}^{b_L}). \end{aligned} \quad \text{(A20)}$$

As for S and T, we use a renormalization scale of order $\mu \sim m_{\rm top}$ to evaluate δg^{b_L}, which errs on the conservative side.

APPENDIX B: CONTRIBUTIONS TO B MESON MIXING

In Sec. IV we estimated the bounds coming from the contributions to $B_{d,s}$ mixing in our model. Here we calculate these contributions in detail. We begin with the simpler case of small bottom Yukawa coupling, based on the bulk masses of Eq. (17), and then we generalize to the large bottom Yukawa case, and show that in fact there are only $\mathcal{O}(1)$ corrections.

1. Small 5D bottom Yukawa

We start with the mass relation of Eq. (16), where for simplicity we omit the overlap correction r_{00}^H (which can be restored at the end) and absorb the $\alpha_{U,D}$ coefficients into the 5D Yukawas. As explained in Sec. IVA, we have to a good approximation $[m_U, Y_U] = 0$. Thus, it is convenient to work in a basis in which Y_U is diagonal. The 5D Yukawa matrices can then be written as

$$Y_U = \lambda_U, \qquad Y_D = V^Q\lambda_D, \quad \text{(B1)}$$

where $\lambda_U = {\rm diag}(Y_u, Y_c, Y_t)$, $\lambda_D = {\rm diag}(Y_d, Y_s, Y_b)$, and V^Q is the misalignment between Y_U and Y_D, or in other words, the 5D equivalent of the CKM matrix.

In order to find the relation between V^Q and $V^{\rm CKM}$, we note that the latter diagonalizes the mass matrix m_D from the left, i.e. it diagonalizes

$$m_D m_D^\dagger \propto F_Q Y_D F_D F_D^\dagger Y_D^\dagger F_Q^\dagger. \quad \text{(B2)}$$

The almost universal F_D's can be thrown away, since we now only care about diagonalization, and for the F_Q's we can pull out a factor of f_{Q^1} ($= f_{Q^2}$), obtaining

$$F_Q \propto {\rm diag}(1, 1, r_Q), \quad \text{(B3)}$$

where we defined[15]

$$r_Q \equiv \frac{f_{Q^3} r_{00}^H(\beta, c_{Q^3}, c_{D^3})}{f_{Q^1} r_{00}^H(\beta, c_{Q^1}, c_{D^3})}. \quad \text{(B4)}$$

Equation (B2) then becomes

$$m_D m_D^\dagger \propto {\rm diag}(1, 1, r_Q)V^Q\lambda_D^2 V^{Q\dagger}{\rm diag}(1, 1, r_Q). \quad \text{(B5)}$$

From this expression, it is simple to find the following relations for the mixing angles[16]

$$V_{12}^Q \sim V_{us}^{\rm CKM}, \qquad V_{13}^Q \sim r_Q V_{td}^{\rm CKM}, \qquad V_{23}^Q \sim r_Q V_{ts}^{\rm CKM}. \quad \text{(B6)}$$

The 5D CKM mixing angles for the third generation are thus larger than the corresponding CKM elements. This is not a surprise, since the hierarchy in the 5D Yukawas is milder than for the masses because of the f_Q's. After diagonalization, we find the mass relations

$$m_{d,s,b} \cong \frac{v}{\sqrt{2}} f_{Q^{1,1,3}} Y_{d,s,b} f_{D^1}, \quad \text{(B7)}$$

where we used the facts that $f_{Q^1} = f_{Q^2}$ in our current realization of the model and that all the f_D's are almost identical.

The matrix D_R which diagonalizes m_D from the right is computed from the following expression:

$$m_D^\dagger m_D \propto F_D^\dagger Y_D^\dagger F_Q^\dagger F_Q Y_D F_D \propto \lambda_D V^{Q\dagger}{\rm diag}(1, 1, r_Q^2)V^Q\lambda_D. \quad \text{(B8)}$$

The resulting mixing angles of D_R are

[15] To be precise, the right-handed bulk mass that is used in the overlap corrections in Eq. (B4) depends on the process in which we use r_Q. Since the largest contributions usually come from the third generation, we defined r_Q with c_{D^3}.

[16] We omit any complex conjugate signs here and below.

$$(D_R)_{12} \sim (r_Q^2 - 1)\frac{Y_d}{Y_s} V_{13}^Q V_{23}^Q \sim (r_Q^2 - 1) r_Q^2 \frac{m_d}{m_s} V_{td}^{\rm CKM} V_{ts}^{\rm CKM},$$
$$(D_R)_{13} \sim \frac{r_Q^2 - 1}{r_Q^2}\frac{Y_d}{Y_b} V_{13}^Q \sim (r_Q^2 - 1)\frac{m_d}{m_b} V_{td}^{\rm CKM}, \tag{B9}$$
$$(D_R)_{23} \sim \frac{r_Q^2 - 1}{r_Q^2}\frac{Y_s}{Y_b} V_{23}^Q \sim (r_Q^2 - 1)\frac{m_s}{m_b} V_{ts}^{\rm CKM},$$

where we used Eqs. (B6) and (B7).

The operators considered in Sec. IV are generated by KK-gluon exchange. The coupling of two zero-mode quarks to a KK-gluon is proportional to $F_Q F_Q^\dagger$ or $F_D^\dagger F_D$ for left- or right-handed couplings, respectively. Applying the appropriate rotation to the down mass basis, the 1–3 entries of the couplings, relevant for B_d mixing, are

$$\begin{aligned}(F_Q F_Q^\dagger)_{13}|_{\rm mass\,basis} \sim{}& V_{tb}^{\rm CKM}(V_{td}^{\rm CKM})^*[f_{Q^3}^2 r_{00}^g(c_{Q^3}) \\ &- f_{Q^1}^2 r_{00}^g(c_{Q^1})],\\ (F_D^\dagger F_D)_{13}|_{\rm mass\,basis} \sim{}& V_{tb}^{\rm CKM}(V_{td}^{\rm CKM})^*\frac{m_d}{m_b}\\ &\times\left[\left(\frac{f_{Q^3} r_{00}^H(\beta, c_{Q^3}, c_{D^3})}{f_{Q^1} r_{00}^H(\beta, c_{Q^1}, c_{D^3})}\right)^2 - 1\right]\\ &\times[f_{D^3}^2 r_{00}^g(c_{D^3}) - f_{D^1}^2 r_{00}^g(c_{D^1})].\end{aligned} \tag{B10}$$

The result for B_s is obtained by the replacements $1 \to 2$ and $d \to s$.

It should be noted that $(D_R)_{12}$ from Eq. (B9) is actually not useful, since when a D_R rotation is applied to $F_D^\dagger F_D$, the 1–2 entry is multiplied by $\sim(f_{D^2}^2 - f_{D^1}^2)$, which is zero in our model. Hence, FCNC processes among the first two generations follow through $(D_R)_{13} \cdot (D_R)_{23}$. This explains why the right-handed current for ϵ_K is suppressed by $r_Q^4 m_d m_s / m_b^2$ relative to the left-handed current, as mentioned in Sec. IV.

2. Large 5D bottom Yukawa

The case where the 5D bottom Yukawa is large is more complicated, but it turns out that it only leads to $\mathcal{O}(1)$ corrections, as we now show. First, Y_U and m_U do not commute anymore, so there is no "natural" basis to adopt. It is therefore useful to define two new matrices, $V^{QD,QU}$, which parametrize the misalignment between $Y_{D,U}$ and C_Q (and equivalently F_Q). Moreover, we need now to compute both D_L and U_L (diagonalizing m_D and m_U from the left, respectively), in order to relate all the above matrices to the CKM matrix. In the following we consider for simplicity only the first relevant terms in the MFV expansion of the 5D spurions.

The first step is to relate V^{QD} and V^{QU} to V^Q. In the basis in which Y_D is diagonal, C_Q from Eq. (14) can be written as

$$C_Q = a_Q \cdot \mathbf{1}_3 + b_D^Q \lambda_D^2 + b_U^Q V^{Q\dagger} \lambda_U^2 V^Q + \cdots, \tag{B11}$$

and it is diagonalized by V^{QD}. We then obtain the following relations:

$$V_{12}^{QD} \sim \frac{V_{13}^Q}{V_{23}^Q} \sim V_{12}^Q,$$
$$V_{13}^{QD} \sim V_{13}^Q\left(\frac{b_U^Q Y_t^2}{b_D^Q Y_b^2 + b_U^Q Y_t^2}\right), \tag{B12}$$
$$V_{23}^{QD} \sim V_{23}^Q\left(\frac{b_U^Q Y_t^2}{b_D^Q Y_b^2 + b_U^Q Y_t^2}\right),$$

where we assumed a specific relation between the V^Q mixing angles for V_{12}^{QD}, which is consistent with the results below (since a similar relation holds for the CKM matrix). Note that the expression in parentheses in the last two mixing angles is of order 1 as long as Y_b is smaller than or comparable to Y_t and $b_{U,D}^Q$ are $\mathcal{O}(1)$. Similarly, in the basis in which Y_U is diagonal, C_Q can be written as

$$C_Q = a_Q \cdot \mathbf{1}_3 + b_U^Q \lambda_U^2 + b_D^Q V^Q \lambda_D^2 V^{Q\dagger} + \cdots. \tag{B13}$$

We then have

$$V_{12}^{QU} \sim V_{12}^Q,$$
$$V_{13}^{QU} \sim V_{13}^Q\left(\frac{b_D^Q Y_b^2}{b_D^Q Y_b^2 + b_U^Q Y_t^2}\right), \tag{B14}$$
$$V_{23}^{QU} \sim V_{23}^Q\left(\frac{b_D^Q Y_b^2}{b_D^Q Y_b^2 + b_U^Q Y_t^2}\right).$$

In this case, if $Y_b < Y_t$ then the expression in parentheses becomes small, and we return to the small bottom Yukawa scenario. We assume that Y_b and Y_t are comparable, so that the expressions in parentheses in Eqs. (B12) and (B14) are $\mathcal{O}(1)$. We then conclude that

$$V^{QD} \sim V^{QU} \sim V^Q. \tag{B15}$$

The next step is to diagonalize from the left the down and up mass matrices, thus expressing D_L and U_L in terms of V^Q. Compared to Eq. (B5) for m_D, we now need to account for the fact that F_D is nonuniversal and not aligned with Y_D. Parametrizing this misalignment by the matrix V^D, Eq. (B5) is generalized to

$$\begin{aligned}m_D m_D^\dagger \propto{}& \mathrm{diag}(1, 1, r_Q) V^Q \lambda_D V^D \mathrm{diag}(1, 1, r_D^2)\\ &\times V^{D\dagger} \lambda_D V^{Q\dagger} \mathrm{diag}(1, 1, r_Q),\end{aligned} \tag{B16}$$

where r_D is defined as r_Q but with $D \leftrightarrow Q$. However, since the leading terms still come from Y_b (so that we can take $Y_d = Y_s = 0$), V^D does not play a role in this diagonalization. Therefore, the result of Eq. (B6) holds also in the current case for the relation between V^Q and D_L, that is

$$V_{12}^Q \sim (D_L)_{12}, \qquad V_{13}^Q \sim r_Q (D_L)_{13}, \qquad V_{23}^Q \sim r_Q (D_L)_{23}. \tag{B17}$$

Applying the same process to m_U, we see that Eq. (B17) also holds after replacing $D_L \to U_L$. Since $V^{\mathrm{CKM}} = U_L D_L^\dagger$, then e.g. for the 2–3 entry,

$$V_{ts}^{\mathrm{CKM}} \sim (U_L)_{22}(D_L)_{23} + (U_L)_{23}(D_L)_{33} \sim \frac{V_{23}^Q}{r_Q}, \tag{B18}$$

where the two terms in the middle are similar in magnitude, but have different phases in general. Thus, we took only one of them as representing the sum (omitting an order 1 correction and the unknown phase). The bottom line is that the relations of Eq. (B6) apply to this case as well, and we have $D_L \sim U_L \sim V^{\mathrm{CKM}}$.

Before continuing, it should be noted that the mass relations in Eq. (B7) are slightly changed to

$$m_{d,s,b} \cong \frac{v}{\sqrt{2}} f_{Q^{1,1,3}} Y_{d,s,b} f_{D^{1,1,3}}, \tag{B19}$$

to include the different c_{D^3}.

In order to estimate D_R, we first need to relate V^D to an already known matrix. In the basis where Y_D is diagonal, V^D diagonalizes C_D, written as

$$C_D = a_D \cdot \mathbf{1}_3 + b_D \lambda_D^2 + d_{DU} \lambda_D V^{Q\dagger} \lambda_U^2 V^Q \lambda_D + \cdots, \tag{B20}$$

considering the relevant leading terms from Eq. (18). The mixing angles of V^D are then given by

$$\begin{aligned} V_{12}^D &\sim \frac{Y_d}{Y_s} Y_t^2 V_{13}^Q V_{23}^Q \left(\frac{d_{DU}}{b_D}\right) \sim r_Q^3 r_D \frac{m_d}{m_s} Y_t^2 V_{td}^{\mathrm{CKM}} V_{ts}^{\mathrm{CKM}}, \\ V_{13}^D &\sim \frac{Y_d}{Y_b} V_{13}^Q \left(\frac{d_{DU} Y_t^2}{b_D + d_{DU} Y_t^2}\right) \sim r_Q^2 r_D \frac{m_d}{m_b} V_{td}^{\mathrm{CKM}}, \\ V_{23}^D &\sim \frac{Y_s}{Y_b} V_{23}^Q \left(\frac{d_{DU} Y_t^2}{b_D + d_{DU} Y_t^2}\right) \sim r_Q^2 r_D \frac{m_s}{m_b} V_{ts}^{\mathrm{CKM}}, \end{aligned} \tag{B21}$$

where we again assume that the expressions in parentheses are $\mathcal{O}(1)$. Now we can generalize Eq. (B8),

$$\begin{aligned} m_D^\dagger m_D \sim{}& \mathrm{diag}(1, 1, r_D) V^{D\dagger} \lambda_D V^{Q\dagger} \mathrm{diag}(1, 1, r_Q^2) \\ &\times V^Q \lambda_D V^D \mathrm{diag}(1, 1, r_D), \end{aligned} \tag{B22}$$

and obtain D_R,

$$\begin{aligned} (D_R)_{12} \sim{}& (r_Q^2 - 1)\frac{Y_d}{Y_s} V_{13}^Q V_{23}^Q \\ &+ V_{12}^D \sim (r_Q^2 - 1) r_Q^2 \frac{m_d}{m_s} V_{td}^{\mathrm{CKM}} V_{ts}^{\mathrm{CKM}} \\ &+ r_Q^3 r_D \frac{m_d}{m_s} Y_t^2 V_{td}^{\mathrm{CKM}} V_{ts}^{\mathrm{CKM}}, \\ (D_R)_{13} \sim{}& \frac{r_Q^2 - 1}{r_Q^2 r_D} \frac{Y_d}{Y_b} V_{13}^Q + \frac{V_{13}^D}{r_D} \sim (r_Q^2 - 1)\frac{m_d}{m_b} V_{td}^{\mathrm{CKM}} \\ &+ r_Q^2 \frac{m_d}{m_b} V_{td}^{\mathrm{CKM}}, \\ (D_R)_{23} \sim{}& \frac{r_Q^2 - 1}{r_Q^2 r_D} \frac{Y_s}{Y_b} V_{23}^Q + \frac{V_{23}^D}{r_D} \sim (r_Q^2 - 1)\frac{m_s}{m_b} V_{ts}^{\mathrm{CKM}} \\ &+ r_Q^2 \frac{m_s}{m_b} V_{ts}^{\mathrm{CKM}}. \end{aligned} \tag{B23}$$

Comparing this to the small bottom Yukawa result in Eq. (B9), we get here for each of the angles the same term plus an additional one, which is of the same order. Since there is a general phase between them, we should just take one of them as the result, so that overall there is an $\mathcal{O}(1)$ correction and an undetermined phase compared to Eq. (B9), as expected. Therefore, we are justified in using Eq. (B10) in our estimates.

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